

# Keldysh approach for non-equilibrium phase transitions in quantum optics: Dicke model in optical cavities

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(Dated: October 16, 2012)

We investigate non-equilibrium phase transitions for driven atomic ensembles interacting with a cavity mode that is subject to dissipation. We first discuss that the relevant modes in the photon spectrum of the driven Dicke model are effectively thermal at low frequencies. The effective temperature of the photons is set by the atom-photon interaction strength and characterizes the static and dynamic critical exponents of the associated superradiance transition. Motivated by these considerations, we develop a general Keldysh path integral approach to describe phase transitions in open quantum-optical systems. This approach is used to compute both the low-frequency and the high frequency dynamics and of the driven Dicke model with Markovian dissipation for the cavity photons and for the atoms. We re-derive the thermal properties of the low-temperature dynamics, while at higher frequencies, we show that the photon and atom distribution functions exhibit quantum and non-equilibrium regimes.

PACS numbers: 42.50.Nn, 03.65.Yz, 05.70.Ln, 31.15.xk

## I. INTRODUCTION

Much interest has recently been directed towards understanding many-body dynamics in open systems away from thermal equilibrium. This subject is not new, as the analogies between threshold phenomena in dynamical systems, such as the laser, and the conventional phase transitions have been recognized over 40 years ago. However, recent experiments with ultracold atoms in optical cavities offer intriguing possibilities to explore the physics of strongly interacting atom-photon systems far away from thermal equilibrium from a new vantage point. Many fundamental concepts of condensed matter physics, ranging from classification of phase transitions to the universal behavior of correlation functions in the vicinity of quantum critical points in the presence of driving and dissipation, need to be revisited in light of these developments.

In this paper, we investigate non-equilibrium phase transitions for driven atomic ensembles interacting with a cavity mode that is subject to dissipation, focusing specifically on the dynamical superradiance transitions and associated self-organization of the atoms observed in Refs. 1–3. Due to the interplay of external driving, Hamiltonian dynamics and dissipative processes, the observed Dicke superradiance transitions [1, 2] exhibit several properties [4–7] which are not present in the closed Dicke model [8–11]. Recently, it was shown that other more interesting quantum many-body phases, such as quantum spin and charge glasses with long-range, random interactions [12–14] mediated by multiple photon modes, could potentially be simulated with many-body cavity QED. Our longer-term objective is to develop a complete theory of these Markovian quantum glasses, including the dissipative processes and glassy phenomenology specific to quantum optics [15]. The present paper establishes the foundation for this

(focussing on the simpler physics with a single cavity mode) en route making a few smaller discoveries that have not been expressed clearly in the literature to date.

In what follows we first review the physics of the non-equilibrium Dicke transition in optical cavities with conventional techniques of quantum optics. Using linearized Heisenberg-Langevin equations, we demonstrate that at low frequencies, close to the phase transition, the system evolves into a thermal state with a high effective temperature, proportional to the atom-photon interaction strength. The static and dynamic critical exponents at the high-temperature phase transitions are analogous to those in a conventional laser (or, more precisely, optical parametric oscillator) threshold.

To treat the interplay between external driving, dissipation and many-body interactions in a more general setting, we next develop a unified approach to describe phase transitions in open quantum systems based on the Keldysh path integral [16–22]. This approach is used to analyze both the low-frequency and the high frequency dynamics of the driven Dicke model with cavity photon loss and atomic spontaneous emission. We re-derive the thermal properties of low-temperature dynamics, while at higher frequencies, we show that the photon and atom distribution functions exhibit quantum and non-equilibrium regimes. We believe the Keldysh calculations are not more involved, and sometimes simpler, than those of the usual quantum optics frameworks [23, 24]. At the same time, they facilitate an easy comparison to other phase transitions of condensed matter physics. One of the objectives of the present paper is to make the Keldysh approach more accessible to the broader quantum optics community. At the same time, we hope that the Keldysh perspective will be helpful for condensed matter physicists to understand driven dissipative atom-photon systems—especially in view of the qualitatively different energy scales and bath properties in quantum optics.

We expect the Keldysh approach to be directly applicable to other dissipative models such as the recently discussed central

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spin model [25] or fermionic lattice models [26–29].

The paper is organized as follows. In Sec. II we introduce the Dicke model and perform a brief analysis with linearized Heisenberg-Langevin equations pointing out that the relevant low-frequency correlations are thermal.

In Sec. III, we map the operators of the master equation and the associated Liouvillians for dissipative processes to the field content of an equivalent real-time, dissipative Keldysh action  $S[\phi, a^*, a]$  with  $\phi$  the atomic fields and  $a^*, a$  the photon field variables. As a warm-up, we first treat the case of a single, decaying cavity mode. Using this simple example, we present in Subsections III A - III D the program that we later in the paper repeat for the photon sector of the Dicke model with cavity photon loss (Sec. IV), and finally also for the atom sector including a simple model for atom decay (Sec. VI).

## II. THERMAL NATURE OF OPEN DICKE TRANSITION

The Dicke Hamiltonian [8, 9] describes  $N$  two-level systems or “qubits” represented by Pauli matrix operators  $\sigma_i^x, \sigma_i^z$  coupled to a quantized photon mode represented by bosonic creation and annihilation operators  $\hat{a}^\dagger, \hat{a}$ :

$$H = \omega_0 \hat{a}^\dagger \hat{a} + \frac{\omega_z}{2} \sum_{i=1}^N \sigma_i^z + \frac{g}{\sqrt{N}} \sum_{i=1}^N \sigma_i^x (\hat{a}^\dagger + \hat{a}). \quad (1)$$

Here  $\omega_0$  is the photon frequency,  $\omega_z$  the level-splitting of the qubits, and  $g$  the qubit-photon coupling, assumed to couple all qubits uniformly to the photon. Eq. (1) is invariant under an Ising-type  $Z_2$ -transformation,  $\hat{a} \rightarrow -\hat{a}$  and  $\sigma_i^x \rightarrow -\sigma_i^x$ . In the thermodynamic limit  $N \rightarrow \infty$ , and for sufficiently strong qubit-photon coupling  $g$ , the ground state of Eq. (1) spontaneously breaks this Ising symmetry and exhibits a phase transition to a “superradiant” phase with a photon condensate  $\langle \hat{a} \rangle$ .

In the context of ultracold dilute gases in optical cavities, Dimer *et al.* [4] proposed to implement the qubits using two hyperfine states of the atoms and showed that, close to the transition, the relevant Hilbert space can be exactly mapped to Eq. (1). Inspired by the work of Dimer *et al.* [4], the qubit states of the Dicke model were realized using two collective motional degrees of the Bose-Einstein condensate of the atoms in the cavity [2, 30], see [31] for a review. In that case  $\omega_z$  becomes a collective recoil frequency and the two Dicke states are components of a dynamically forming charge density wave. This open realization of the Dicke model in optical cavities with pumped atoms is different from the closed system Dicke model Eq. (1) for three principal reasons, which we now explain one-by-one:

**1. Coherent drive:** The mapping to the Dicke model is performed in a frame that rotates with the (optical) pump frequency of the driving laser. In this rotating frame, the explicit time-dependence of the original Jaynes-Cummings type Hamiltonian [54] is “gauged” away. For the effective parameters of the Dicke model Eq. (1), this leads to a shift of the cavity and atomic frequency scales: in particular  $\omega_0 = \omega_c - \omega_p$  with  $\omega_c$  the bare cavity frequency and  $\omega_p$  the pump frequency (cf. Appendix A). This “hidden” time-dependence of the ef-

fective parameters of the Dicke model turns the electromagnetic vacuum surrounding the cavity into a *Markovian non-equilibrium bath*.

**2. Cavity photon loss:** In addition to the coherent dynamics generated by the Hamiltonian Eq. (1), there is a dissipative contribution consisting of cavity loss and spontaneous emission for the atoms. The complete dynamics of the system density matrix is modeled by a Markovian master equation,

$$\partial_t \rho = -i[H, \rho] + \sum_{\alpha} \kappa_{\alpha} (2L_{\alpha} \rho L_{\alpha}^{\dagger} - \{L_{\alpha}^{\dagger} L_{\alpha}, \rho\}), \quad (2)$$

where the curly brackets  $\{, \}$  denote the anti-commutator and  $L_{\alpha}$  is a set of Lindblad or quantum jump operators, where  $\alpha$  labels different degrees of freedom as well as different positions. The second piece of (2) describes the dissipative dynamics in terms of a Liouville operator  $\mathcal{L}$  in Lindblad form, acting on the density matrix  $\rho$ . Here we have two contributions: cavity photon loss and atomic spontaneous emission. The former is modeled by the Liouvillian:

$$\mathcal{L}_{\text{cav}} \rho = \kappa (2\hat{a} \rho \hat{a}^{\dagger} - \{\hat{a}^{\dagger} \hat{a}, \rho\}), \quad (3)$$

where  $\kappa$  is an effective decay rate (inverse lifetime) of a cavity photon of the order a few MHz [2].

**3. Atomic spontaneous emission:** Modelling the dissipative dynamics of the atoms depends on the specific implementation of the driven Dicke model usually involving local processes of each two-level atom separately. In Sect. VI, we account for spontaneous emission in an approximate way by resorting to a simplified effective low frequency model.

### A. Heisenberg-Langevin analysis

In this section, we study the effective dynamics of the Dicke model near the phase transition using conventional quantum optical techniques, namely, the Heisenberg equations of motion for the atomic and photon operators. We begin by neglecting spontaneous emission of the atoms, but include cavity decay into a Markovian reservoir. In this case the dynamics is described by a set of operator equations:

$$\dot{\hat{a}} = -i\omega_0 \hat{a} - \kappa \hat{a} - \frac{ig}{\sqrt{N}} \sigma_i^x + \mathcal{F}, \quad (4)$$

$$\dot{\sigma}_i^+ = i\omega_z \sigma_i^+ - \frac{ig}{\sqrt{N}} \sigma_i^z (\hat{a} + \hat{a}^\dagger), \quad (5)$$

$$\dot{\sigma}_i^z = -2 \frac{ig}{\sqrt{N}} (\sigma_i^+ - \sigma_i^-) (\hat{a} + \hat{a}^\dagger). \quad (6)$$

Here the force  $\mathcal{F} = \mathcal{F}(t)$  is a stochastic operator satisfying  $\langle \mathcal{F}(t) \mathcal{F}^\dagger(t') \rangle = 2\kappa \delta(t - t')$  and  $\langle \mathcal{F}(t)^\dagger \mathcal{F}(t) \rangle = 0$ . This term is needed in order to preserve the commutation relation  $[\hat{a}(t), \hat{a}^\dagger(t)] = 1$ , which would otherwise exponentially decrease. To analyze the dynamics below superradiance threshold, we assume that the atoms are fully polarized  $S^z = \frac{1}{2} \sum_i \sigma_i^z \approx -N/2$ , and neglect non-linear terms in the equations of motion. The resulting operator equa-

tions can be solved exactly in the Fourier domain as has been done in detail in the work of Dimer *et al.* [4] and we will not repeat their calculations here. We just add one simple point to their comprehensive analysis, namely that the relevant low-frequency dynamics of the photons occurs in the presence of a finite effective temperature. This is most easily seen from the equation of motion for the real-valued phase-space coordinate  $x(\omega) = (\hat{a}(\omega) + \hat{a}^\dagger(-\omega))/\sqrt{2\omega_0}$ . Defining stochastic force operators  $f(\omega) = \frac{1}{\sqrt{2\omega_0}} [\mathcal{F}(\omega)(\kappa - i(\omega_0 + \omega)) + \mathcal{F}^\dagger(-\omega)(\kappa + i(\omega_0 - \omega))]$ , we can write the equation of motion as

$$\left( (-\kappa + i\omega)^2 + \omega_0^2 - \frac{4g^2\omega_z\omega_0}{\omega_z^2 - \omega^2} \right) x(\omega) = f(\omega). \quad (7)$$

The force operator satisfies

$$\frac{1}{2} \langle f(\omega)f(\omega') + f(\omega')f(\omega) \rangle = \kappa \frac{\kappa^2 + \omega_0^2 + \omega^2}{\omega_0} \delta(\omega + \omega'). \quad (8)$$

At low frequencies, we can neglect high-order terms in  $\omega$ . Eq. (7) becomes identical to the Langevin equation of a classical particle in a harmonic potential with oscillation frequency  $\alpha$ , defined by

$$\alpha^2 = \kappa^2 + \omega_0^2 - \frac{4g^2\omega_0}{\omega_z}, \quad (9)$$

and friction constant  $2\kappa$ . This oscillator becomes dynamically unstable at the Dicke transition, where the frequency  $\alpha$  vanishes, given by the critical coupling

$$g_c = \sqrt{\frac{\omega_0^2 + \kappa^2}{4\omega_0}} \omega_z. \quad (10)$$

The correlation function of the stochastic force operators on the right-hand-side of Eq. (8) becomes identical to the “noise correlations” provided by an equilibrium bath at a finite temperature

$$T_x^{\text{eff}} = \frac{\omega_0^2 + \kappa^2}{4\omega_0} = \frac{g_c^2}{\omega_z}. \quad (11)$$

However, in contrast to the temperature in an equilibrium problem, the low-frequency effective temperature here is not a global property of the system, but is in general observable dependent. In Sec. IV D, we present a systematic, generalizable way to extract low-frequency effective temperatures based on observable-dependent fluctuation-dissipation relations. We already here quote the effective temperature for the atoms (see Sec. VI for the details of the computation):

$$T_\phi^{\text{eff}} = \frac{\gamma^2 + \omega_z^2}{4\omega_z}, \quad (12)$$

where  $\omega_z$  is the recoil energy and  $\gamma$  an effective single atom decay rate. Because  $T_\phi^{\text{eff}} \neq T_x^{\text{eff}}$ , the different parts of the driven system thus do not equilibrate to each other and, although the

dominant low-frequency correlations are thermal, the system is not in a global thermal state. We remark that our definition of effective temperature does not coincide with those commonly used in laser theory, as discussed in Sec. IV D: the former relates the fluctuations of the field to its response in the rotating frame, while the latter compares the fluctuations of the field to an equilibrium situation in the lab frame.

We now demonstrate that the key critical exponent characterizing the non-equilibrium Dicke transition is dominated by the thermal modes in the spectrum. To that end, we compute the “photon flux exponent” [5] for photon occupation in the vicinity of the transition. The photon occupation is

$$\begin{aligned} 2\langle n \rangle + 1 &= \langle \hat{a}^\dagger \hat{a} + \hat{a} \hat{a}^\dagger \rangle = 2\omega_0 (\langle x^2 \rangle + \langle p^2 \rangle) \\ &= 2\omega_0 \left( 1 + \frac{\kappa^2}{\omega_0^2} \right) \langle x^2 \rangle, \end{aligned} \quad (13)$$

where the expression for  $p = i(\hat{a} - \hat{a}^\dagger)/\sqrt{2\omega_0}$  can be obtained in complete analogy to that of  $x$ . We can compute  $\langle x^2 \rangle$  using the equilibrium partition function [32, 33]  $Z = \exp(-F/T_{aa}^{\text{eff}})$ , with  $F = \frac{1}{2}\alpha^2 x^2$ . Performing the Gaussian integral yields

$$2\langle n \rangle + 1 = 2 \left( 1 + \frac{\kappa^2}{\omega_0^2} \right) \omega_0 \frac{T_{aa}^{\text{eff}}}{\alpha^2} \sim \frac{1}{|g - g_c|^{\nu_x}}, \quad (14)$$

with the static photon flux exponent  $\nu_x = 1$ , as already noted in Ref. 5, 6. There the authors also noted the discrepancy with the zero temperature case [11], where  $\langle n \rangle \sim 1/\alpha \sim 1/|g^2 - g_c^2|^{1/2}$  and  $\nu_x = 1/2$ . Here we offer an explanation of this discrepancy: the dominant photon modes at the Dicke transition have a non-zero effective temperature.

The above results indicate that, from the point of view of phase transitions, it is incorrect to call the driven Dicke transition a quantum phase transition even though it is “made of quantum ingredients” (two collective motional states of a BEC [2]), and the photon spectrum contains a squeezed quantum mode. Instead, it should be regarded as a classical phase transition governed by the properties of the dynamical universality class of the classical Ising model, the so-called “Model A” of Hohenberg and Halperin [34]. The effect that dissipation induces a finite effective temperature is not new. In several other condensed matter systems [35–37], the coupling to a non-equilibrium bath typically admixes the pure many-body states of the closed system, transforming pure quantum phase transitions into thermal phase transitions (see also [17, 18, 25, 38, 39]). What is perhaps more surprising is that neither the low-frequency effective temperature for the photons nor for the atoms is set by the cavity loss rate  $\kappa$ , but instead set by the atom-photon interaction  $g$  and the effective single atom parameters, respectively.

Being complex-valued objects, the photons have two normal modes. As we saw, one quadrature is thermally amplified and diverges at the Dicke transition. Its orthogonal quadrature is quantum squeezed [4], remains gapped at the transition and therefore does not influence the thermal nature of the phase transition. These attenuated and amplified quadratures arise naturally as the eigenmodes of the photon correlation func-

tion (see Sec. IV B).

The remainder of the paper is dedicated to the development of a unified, and generalizable, Keldysh approach. The Dicke model will be used as a prototypical test object and we compare our results to those of other approaches, where available.

### III. KELDYSH APPROACH FOR CAVITY VACUUM

In this section, we introduce the real-time Keldysh formalism and fix our notation by considering the case of a single electromagnetic mode in an open cavity (without atoms). We also indicate how to relate the operators of the master equation (2) to the field content and choice of time contour in a Keldysh action (see also Refs. [16, 19, 40]).

The decay of a single boson (the cavity photon) into a continuum of modes (the external vacuum) is described by the master equation

$$\partial_t \rho = -i[\omega_0 \hat{a}^\dagger \hat{a}, \rho] + \kappa(2\hat{a}\rho\hat{a}^\dagger - \{\hat{a}^\dagger \hat{a}, \rho\}). \quad (15)$$

This equation results from the (fully unitary) Heisenberg equation for the coupled system-bath setting, where the system is described by the degrees of freedom  $\hat{a}, \hat{a}^\dagger$  and the bath by a continuum of harmonic modes. Eq. (15) is obtained by eliminating (“integrating-out”) the bath in the Born-Markov and rotating wave approximation (cf. e.g. [23]). The integration of the bath variables gives rise to an effective evolution including dissipative terms. Performing the same program in path integral formulation, we obtain the *Markovian dissipative action*

$$S_a = \int_{-\infty}^{\infty} dt \left( a_+^*(i\partial_t - \omega_0)a_+ - a_-^*(i\partial_t - \omega_0)a_- - i\kappa[2a_+a_-^* - (a_+^*a_+ + a_-^*a_-)] \right). \quad (16)$$

In the path integral formalism, the quantum mechanical operators are replaced by fluctuating, time-dependent and complex-valued fields (we omit the time argument for notational simplicity). The fact that the density matrix can be acted on from both sides, as reflected in the Heisenberg commutator structure of the original evolution equation, finds its counterpart in the presence of a forward (+) and backward (-) component of the fields. The former is associated to an action on the density matrix from the left, and the latter to the right. Indeed, in the first of line of Eq. (16) there is a relative minus sign between the terms involving the two components, reflecting the Heisenberg commutator structure of Eq. (2). The terms in the second line instead display the characteristic Lindblad form; the “jump” or “recycling” term is represented by an explicit coupling of the two contours.

It is convenient [20–22] to introduce “center-of-mass” and “relative” field coordinates,  $a_{cl} = (a_+ + a_-)/\sqrt{2}$ ,  $a_q = (a_+ - a_-)/\sqrt{2}$ . These new coordinates are often referred to as “classical” and “quantum” fields, because the first can acquire an expectation value while the second one cannot in the absence of sources. In this basis, and going to frequency space, we

write

$$S_a = \int_{\omega} (a_{cl}^*, a_q^*) \begin{pmatrix} 0 & [G^A]^{-1}(\omega) \\ [G^R]^{-1}(\omega) & D^K(\omega) \end{pmatrix} \begin{pmatrix} a_{cl} \\ a_q \end{pmatrix}, \quad (17)$$

where we used the notation  $\int_{\omega} = \int_{-\infty}^{\infty} \frac{d\omega}{2\pi}$ , and  $a_{cl,q}(t) = \int_{\omega} e^{-i\omega t} a_{cl,q}(\omega)$ . This classical-quantum basis is often referred to as RAK basis: the entries are the inverse Retarded (lower left) and Advanced (upper right) Green’s functions, and the inverse Keldysh component. The RAK action Eq. (17) can be easily inverted to deliver the photonic Green’s functions

$$\begin{pmatrix} G^K(\omega) & G^R(\omega) \\ G^A(\omega) & 0 \end{pmatrix} = \begin{pmatrix} 0 & [G^A]^{-1}(\omega) \\ [G^R]^{-1}(\omega) & D^K(\omega) \end{pmatrix}^{-1}, \quad (18)$$

where the Keldysh Green’s functions is a matrix product

$$G^K(\omega) = -G^R(\omega)D^K(\omega)G^A(\omega). \quad (19)$$

For the open cavity of Eq. (16) the RAK inverse Green’s functions are

$$[G^{-1}]^{R/A} = \omega - \omega_0 + \Sigma_a^{R/A}, \quad D^K = \Sigma_a^K, \quad (20)$$

with the “self-energies”

$$\Sigma_a^A = -i\kappa, \quad \Sigma_a^R = +i\kappa, \quad \Sigma_a^K = 2i\kappa. \quad (21)$$

It is a key property of a Markovian system that the Keldysh component  $\Sigma_a^K$  in Eq. (21) is frequency independent. As can be seen from App. A, this is due to a separation of scales between (i) the large pump ( $\omega_p$ ) and cavity frequency ( $\omega_c$ ), both optical frequencies in the Tera Hertz range ( $\sim 10^{14}$  Hz corresponding to temperatures  $T \sim 10^4 K$ ), and (ii) the characteristic frequencies of the electromagnetic vacuum outside the cavity ( $\lesssim 10^{12}$  Hz corresponding to temperatures  $T \lesssim 300 K$ ).

In the literature, it is often argued that a frequency independent, nonzero inverse Keldysh component indicates an effective finite temperature state. The Markovian lossy cavity is a simple counterexample: Even though the inverse Keldysh component is constant  $\sim 2i\kappa$ , the state is pure and the effective temperature zero, as we will argue below.

We next introduce the key propagators that encode the systems’ *response* and *correlation* functions. In equilibrium, the two are rigidly related by the Bose (or Fermi) distribution function; out-of-equilibrium, no such *a priori* knowledge is available, and it is important to distinguish them.

#### A. Cavity spectral response function

The *spectral response function* encodes the system’s response to active, external perturbations such as time-modulated external fields coupling to spin or charge operators, for example. The spectral response function is the difference

between the retarded and advanced Green's function:

$$\mathcal{A}(\omega) = i(G^R(\omega) - G^A(\omega)). \quad (22)$$

In the scalar case considered here, we have

$$\mathcal{A}_{aa^\dagger}(\omega) = -2\text{Im}G^R(\omega), \quad (23)$$

The frequency-integrated spectral response function is normalized to unity, because of the exact commutator relation of the bosonic degrees of freedom

$$\int \frac{d\omega}{2\pi} \mathcal{A}_{aa^\dagger}(\omega) = \langle [\hat{a}, \hat{a}^\dagger] \rangle = 1. \quad (24)$$

This “sum rule” is an exact property of the theory valid in- and out-of-equilibrium. In our example of one cavity mode,

$$\mathcal{A}_{aa^\dagger}(\omega) = \frac{2\kappa}{(\omega - \omega_0)^2 + \kappa^2}. \quad (25)$$

### B. Cavity correlation function

The *correlation function* encodes the system's internal correlations, for example the frequency-resolved photon spectrum of the intracavity photon fields. In the steady state, the photon correlation function is related to the Keldysh Green's function by

$$C_{aa^\dagger}(t) = \langle \{\hat{a}(t), \hat{a}^\dagger(0)\} \rangle = \langle \hat{a}(t)\hat{a}^\dagger(0) + \hat{a}^\dagger(0)\hat{a}(t) \rangle = iG^K(t). \quad (26)$$

Here the last identity is valid only in the specific case of a scalar Keldysh Green's function. At equal times this relation results in

$$C_{aa^\dagger}(0) = 2\langle \hat{a}^\dagger \hat{a} \rangle + 1 = iG^K(t=0) = i \int \frac{d\omega}{2\pi} G^K(\omega). \quad (27)$$

For the single decaying cavity mode, characterized by Eqs. (17-21), it is easy to show that  $C_{aa^\dagger}(\omega) = \mathcal{A}_{aa^\dagger}(\omega)$  and the frequency integral over the Keldysh Green's function is unity yielding  $\langle \hat{a}^\dagger \hat{a} \rangle = 0$ . As expected, the steady state corresponds to the cavity vacuum.

### C. Comparison with a closed system at equilibrium

In the absence of dissipation, the cavity becomes an isolated harmonic oscillator. Its inverse Green's functions are still given by (17) with the self-energies serving only as regularization parameters. At equilibrium,

$$\begin{aligned} \Sigma_{a,\text{EQ}}^A &= -i\epsilon, & \Sigma_{a,\text{EQ}}^R &= +i\epsilon, \\ \Sigma_{a,\text{EQ}}^K(\omega) &= 2i\epsilon \coth \left[ \frac{\omega}{2T} \right]. \end{aligned} \quad (28)$$

Here  $\epsilon \rightarrow 0$  at the end of the calculation and  $T$  is the actual temperature. Note that the Keldysh component is odd with

respect to the frequency  $\Sigma_{a,\text{EQ}}^K(-\omega) = -\Sigma_{a,\text{EQ}}^K(\omega)$ , while in the Markovian system (21) it is even.

Using Eq. (28), we obtain

$$\begin{aligned} \mathcal{A}_{aa^\dagger}^{\text{EQ}}(\omega) &= \lim_{\epsilon \rightarrow 0} \frac{2\epsilon}{(\omega - \omega_0)^2 + \epsilon^2} = 2\pi\delta(\omega - \omega_0), \\ C_{aa^\dagger}^{\text{EQ}}(\omega) &= \lim_{\epsilon \rightarrow 0} \frac{2\epsilon}{(\omega - \omega_0)^2 + \epsilon^2} \coth \frac{\omega}{2T} = 2\pi \coth \frac{\omega}{2T} \delta(\omega - \omega_0), \end{aligned} \quad (29)$$

In this noninteracting case, the spectral response function is fully centered at the isolated mode with frequency  $\omega_0$ . We observe that formally, the thermodynamic equilibrium limit can be seen as a situation with an infinitesimal loss, replacing  $\kappa \rightarrow \epsilon$ , and the replacement in the inverse Keldysh component  $2i\kappa \rightarrow 2i\epsilon \coth \frac{\omega}{2T}$ .

### D. Cavity distribution function and low-frequency effective temperature

The response and correlations allow us to define a fluctuation-dissipation relation, by introducing the *distribution function*  $F(\omega)$ :

$$\begin{aligned} G^K(\omega) &= G^R(\omega)F(\omega) - F(\omega)G^A(\omega) \\ \Leftrightarrow D^K(\omega) &= [G^R(\omega)]^{-1}F(\omega) - F(\omega)[G^A(\omega)]^{-1}, \end{aligned} \quad (30)$$

where the equivalence holds due to Eq. (19). At thermal equilibrium the distribution  $F$  is universal and equals to the unit matrix times

$$\begin{aligned} F^{\text{EQ}}(\omega) &= \coth \frac{\omega}{2T} = 2n_B(\frac{\omega}{T}) + 1, \\ F_{T=0}^{\text{EQ}}(\omega) &= \text{sign}(\omega), \\ F_{\omega \ll T}^{\text{EQ}}(\omega) &\approx \frac{2T}{\omega} + \dots \end{aligned} \quad (31)$$

with the Bose distribution  $n_B(x) = (\exp x - 1)^{-1}$ . The unit matrix in field space signals detailed balance between all subparts of the system.

In the present case, the system is out-of-equilibrium due to its driven and dissipative nature. A notion of a temperature is not *a priori* meaningful: Neither must the driven system equilibrate to an external heat bath with temperature  $T$  (in our case, due to the separation of scales underlying the Markov approximation, this temperature would be effectively  $T = 0$  compared to the system scales), nor do the different subparts of the system have to equilibrate with respect to each other. In this work, we argue that a notion of a temperature nevertheless emerges as a universal feature of the low frequency domain of Markovian systems. It is introduced by computing the  $F$  matrix through Eq. (30) and comparing the low-frequency behavior of its eigenvalues with the equilibrium result of Eq. (31). In particular, if  $F$  has a thermal infrared enhancement  $\sim 1/\omega$  for small frequencies, its dimensionful coefficient is identified as an effective temperature. This notion of a “low frequency effective temperature” (LET) becomes particularly relevant in the vicinity of a phase transition, where the spectral weight encoded in  $G^R, G^A$  is concentrated near zero fre-

quency. Below, we will use this concept to establish a connection of Markovian quantum systems to the classical theory of dynamical universality classes according to Hohenberg and Halperin [34]. Moreover, we find that, while all governed by the  $1/\omega$  divergence in the distribution function, different subparts of the system exhibit different LETs. In contrast to the global temperature present in thermodynamic equilibrium, the LET is not an external parameter but rather a system immanent quantity, determined by the interplay of unitary and dissipative dynamics.

In case of a decaying cavity, the Green's functions are scalars and we can easily invert (30) to obtain

$$F(\omega) = \frac{G^K(\omega)}{G^R(\omega) - G^A(\omega)} = \frac{C_{aa^\dagger}(\omega)}{\mathcal{A}_{aa^\dagger}(\omega)} = 1. \quad (32)$$

confirming that the cavity vacuum has a zero effective temperature. Moreover, as for a pure quantum state in the equilibrium case at  $T = 0$ , the distribution function here also squares to a unit matrix,  $F^2(\omega) = 1$ .

An important difference between the zero temperature equilibrium case and Markovian case appears in the sign of the distribution function. In the former case (as for any equilibrium distribution),  $F(\omega)$  is anti-symmetric with respect to the frequency  $\omega$ . On the contrary, for a Markovian bath the distribution function is symmetric with respect to  $\omega$ ; one signature of a strongly-out-of-equilibrium system.

#### IV. KELDYSH APPROACH FOR PHOTON OBSERVABLES

We now analyze the Dicke model Eq. (1) with the path integral approach explained in the previous section. We include cavity photon loss but defer the inclusion of spontaneous emission to Sec. VI. Assuming homogeneous qubit-cavity coupling, one can use the large- $N$  strategy of Emary and Brandes [11]. We introduce collective large- $N$  spin operators  $S_z = \frac{1}{2} \sum_{i=1}^N \sigma_i^z$  and  $S^\pm = \frac{1}{2} \sum_{i=1}^N (\sigma_i^\pm + \sigma_i^\mp)$ , to write the Dicke model (1) in terms of one large spin coupled to the cavity photon mode,

$$H = \omega_0 \hat{a}^\dagger \hat{a} + \omega_z S_z + \frac{2g}{\sqrt{N}} S^x (\hat{a}^\dagger + \hat{a}). \quad (33)$$

We then express the spin in terms of a Holstein-Primakoff [11, 41] boson operator  $\hat{b}$ , defined by  $S_z = -N/2 + \hat{b}^\dagger \hat{b}$ ,  $S^+ = \sqrt{N - \hat{n}} \hat{b}^\dagger \approx \sqrt{N} (1 - \hat{n}/(2N)) \hat{b}^\dagger$ , and  $S^- = (S^+ + S^-)/2$ . Neglecting unimportant constants we obtain the normal ordered Hamiltonian:

$$H = \omega_0 \hat{a}^\dagger \hat{a} + \omega_z \hat{b}^\dagger \hat{b} + g (\hat{a} + \hat{a}^\dagger) (\hat{b} + \hat{b}^\dagger - \frac{1}{2N} \hat{b}^\dagger (\hat{b}^\dagger + \hat{b}) \hat{b}). \quad (34)$$

At  $N \rightarrow \infty$ , the last, non-quadratic term vanishes and the problem reduces to a linear system of two coupled bosonic degrees of freedom one of which (the cavity mode  $\hat{a}$ ) decays into a Markovian bath. As outlined in section III, we can transform the Liouvillian Eq. (3), with Hamiltonian Eq. (34), into

an equivalent Keldysh action with

$$S = S_a + S_b + S_{ab} \quad (35)$$

$$S_b = \int_{\omega} (b_{cl}^*, b_q^*) \begin{pmatrix} 0 & \omega - \omega_z \\ \omega - \omega_z & 0 \end{pmatrix} \begin{pmatrix} b_{cl} \\ b_q \end{pmatrix}$$

where  $S_a$  is given by Eq. (17) and the interaction in terms of the “classical” and “quantum” fields reads

$$S_{ab} = -g \int_{\omega} (1 - \frac{1}{2N}) [(a_q + a_q^*)(b_{cl} + b_{cl}^*) + (a_{cl} + a_{cl}^*)(b_q + b_q^*)] \\ - \frac{1}{N} \left\{ [(a_{cl} + a_{cl}^*)(b_q + b_q^*) + (a_q + a_q^*)(b_{cl} + b_{cl}^*)] [b_{cl}^* b_{cl} + b_q^* b_q] \right. \\ \left. + [(a_{cl} + a_{cl}^*)(b_{cl} + b_{cl}^*) + (a_q + a_q^*)(b_q + b_q^*)] [b_{cl}^* b_q + b_q^* b_{cl}] \right\}.$$

We now demonstrate that the static saddle point solutions of this action reproduce the results of other approaches [4]. Varying  $S$  with respect to the quantum components of the fields and substituting  $a_{cl}(t) = a_0$ ,  $b_{cl}(t) = b_0$ ,  $a_q = 0$ ,  $b_q = 0$ , we obtain the coupled equations

$$\frac{\partial S}{\partial a_q^*} = (-\omega_0 + i\kappa) a_0 - g \left( 1 - \frac{1}{N} b_0^2 \right) 2b_0 = 0, \quad (36)$$

$$\frac{\partial S}{\partial b_q^*} = -\omega_z b_0 - g \left( 1 - \frac{3}{N} b_0^2 \right) (a_0 + a_0^*) = 0. \quad (37)$$

where we chose  $b_0 = b_0^*$ . As expected, these saddle-point equations admit solutions with non-zero “ferromagnetic” moment  $b_0$ , and super-radiant photon condensate  $a_0$ ,

$$b_0 = \pm \frac{1}{2} \sqrt{N} \sqrt{\frac{g^2 - g_c^2}{g^2}} \quad a_0 = \pm \sqrt{N} \frac{\sqrt{g^2 - g_c^2}}{\omega_0 - i\kappa} \quad (38)$$

for atom-photon couplings larger than a critical value

$$g > g_c = \sqrt{\frac{\omega_0^2 + \kappa^2}{4\omega_0}} \omega_z. \quad (39)$$

We are now going to integrate-out the atomic field  $b$  and obtain an effective action describing the photons in the normal phase ( $g < g_c$ ), where  $a_0 = b_0 = 0$ . In App. B, we give the corresponding expressions in the superradiant phase.

In the limit of  $N \rightarrow \infty$ , we can safely neglect the terms proportional to  $1/N$  in (35). We can re-write Eq. (35) as  $8 \times 8$  matrix multiplying the 8-component fields:

$$V_8(\omega) = \begin{pmatrix} a_{cl}(\omega) \\ a_{cl}^*(-\omega) \\ b_{cl}(\omega) \\ b_{cl}^*(-\omega) \\ a_q(\omega) \\ a_q^*(-\omega) \\ b_q(\omega) \\ b_q^*(-\omega) \end{pmatrix} \quad (40)$$

with the action

$$S_N = \frac{1}{2} \int_{\omega} V_8^{\dagger}(\omega) \begin{pmatrix} 0 & [G_{4 \times 4}^A]^{-1}(\omega) \\ [G_{4 \times 4}^R]^{-1}(\omega) & D_{4 \times 4}^K \end{pmatrix} V_8(\omega). \quad (41)$$

The subscript N stands for normal phase and the dagger  $\dagger$  denotes transposition and complex conjugation. The block entries are  $4 \times 4$  Green's functions given by

$$[G_{4 \times 4}^R]^{-1}(\omega) = \begin{pmatrix} \omega - \omega_0 + i\kappa & 0 & -g & -g \\ 0 & -\omega - \omega_0 - i\kappa & -g & -g \\ -g & -g & \omega - \omega_z & 0 \\ -g & -g & 0 & -\omega - \omega_z \end{pmatrix} \quad (42)$$

$$D_{4 \times 4}^K = 2i \text{diag}(\kappa, \kappa, 0, 0).$$

To obtain the photon-only action, we now integrate out the Holstein-Primakoff field  $b$  and get a Keldysh functional integral that goes only over the photon fields:  $Z^K = \int D[a^*, a] e^{iS_{\text{photon}}[a^*, a]}$ , with the photon-only action

$$S_{\text{photon}}[a^*, a] = \int_{\omega} A_4^{\dagger}(\omega) \begin{pmatrix} 0 & [G_{2 \times 2}^A]^{-1}(\omega) \\ [G_{2 \times 2}^R]^{-1}(\omega) & D_{2 \times 2}^K(\omega) \end{pmatrix} A_4(\omega). \quad (43)$$

The photon four-vector collects the classical and quantum field components

$$A_4(\omega) = \begin{pmatrix} a_{cl}(\omega) \\ a_{cl}^*(-\omega) \\ a_q(\omega) \\ a_q^*(-\omega) \end{pmatrix}. \quad (44)$$

and the block entries are  $2 \times 2$  photon Green's functions which we now analyze one-by-one.

### A. Photon spectral response function

The inverse retarded Green's function of the photons is

$$[G_{2 \times 2}^R]^{-1}(\omega) = \begin{pmatrix} \omega - \omega_0 + i\kappa + \Sigma^R(\omega) & \Sigma^R(\omega) \\ [\Sigma^R(-\omega)]^* & -\omega - \omega_0 - i\kappa + [\Sigma^R(-\omega)]^* \end{pmatrix}, \quad (45)$$

where the interaction induced photon-self energy reads

$$\Sigma^R(\omega) = -\frac{2g^2\omega_z}{\omega^2 - \omega_z^2}. \quad (46)$$

The characteristic frequencies of the system are defined by the zeros of the determinant  $[G_{2 \times 2}^R]^{-1}(\omega)$ , corresponding to the poles of the response function  $G_{2 \times 2}^R(\omega)$ . Due to the symmetry

$$\sigma_x [G_{2 \times 2}^R(-\omega)]^* \sigma_x = G_{2 \times 2}^R(\omega), \quad (47)$$

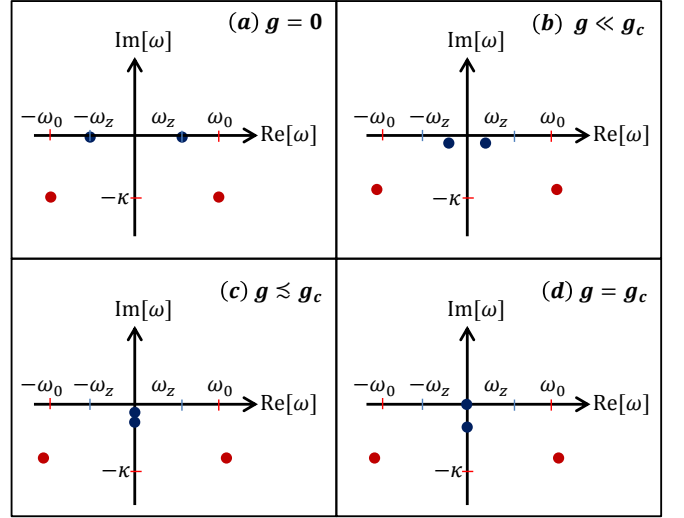


FIG. 1: (Color online) Schematic plot of the position of the poles of the retarded Green's function Eq. (48). (a) At zero coupling  $g = 0$ , two poles can be associated with the photonic mode  $\omega = \pm\omega_0 + i\kappa$  and two with the atomic mode  $\omega = \pm\omega_z$ . (b) In the presence of a finite coupling  $g$ , the modes hybridize and the corresponding frequencies are shifted in opposite directions. (c) When approaching the transition, two solutions become purely imaginary and correspond to damped modes. (d) At the transition point  $g = g_c$  one of the poles approaches zero, making the system dynamically unstable.

the poles come in pairs, such that  $\{\lambda\} = \{-\lambda^*\}$ , meaning that they either are pure imaginary or come in pairs with opposite real part. The explicit solution of

$$0 = \frac{1}{\det[G_{2 \times 2}^R(\omega)]} = \left( \omega_0 + \frac{2g^2\omega_z}{\omega^2 - \omega_z^2} \right)^2 - (\omega + i\kappa)^2 - \left( \frac{2g^2\omega_z}{\omega^2 - \omega_z^2} \right)^2 \quad (48)$$

yields four poles, schematically plotted in Fig. 1. Note that, in the vicinity of the phase transition, two poles become purely imaginary. This phenomenon seems to apply to generic dissipative transitions as we further describe in App. C. Indeed it has been previously observed for a dissipative critical central spin model [25]. Overdamping of collective modes, due to a similar mechanism, has also been found in dissipative multi-mode systems in symmetry broken phases [17, 18].

The imaginary part of the first diagonal element of  $G_{2 \times 2}^R(\omega)$  corresponds to the photon spectral response function  $\mathcal{A}_{aa^*}$ , defined in (23), and is plotted in Fig. 2 for different values of the coupling  $g$ . In the absence of atom-photon coupling  $g = 0$  (dotted curve) there is a single resonance peak at frequency  $\omega = \omega_0$ , broadened by the cavity decay rate  $\kappa$ . A finite coupling  $g$  (dashed curve) “collectively Rabi-splits” the resonance [42] in two distinct peaks, corresponding to two distinct poles of the system. Upon approaching the Dicke transition (solid curve), the spectral weight is shifted towards the low-frequency pole; a precursor to the superradiant cavity mode.

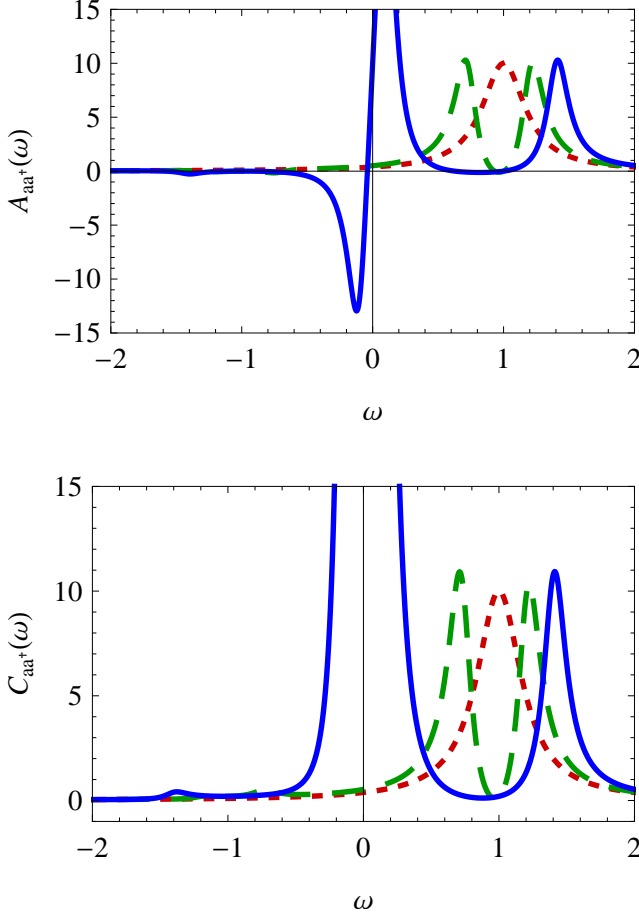


FIG. 2: (Color online) Photon spectral response function  $A_{aa^\dagger}(\omega)$  and correlation function  $C_{aa^\dagger}(\omega)$  as a function of real frequencies  $\omega$ . Numerical parameters:  $\omega_0 = \omega_z = 1$ ,  $\kappa = 0.2$  (leading to  $g_c \approx 0.51$ ), and  $g=0$  (dotted), 0.25 (dashed), 0.5 (solid).

### B. Photon correlation function

The inverse Keldysh component of the action Eq. (43) is

$$D_{2 \times 2}^K = \begin{pmatrix} 2i\kappa & 0 \\ 0 & 2i\kappa \end{pmatrix} \quad (49)$$

and the Keldysh Green's function  $G_{2 \times 2}^K(\omega) = -G_{2 \times 2}^R(\omega)D_{2 \times 2}^K G_{2 \times 2}^A(\omega)$  is a  $2 \times 2$  matrix.

The first diagonal element of  $G_{2 \times 2}^K$ ,  $C_{aa^\dagger} = i(1 \ 0)G_{2 \times 2}^K(1 \ 0)^T$  corresponds to the photon correlation function defined in Eq. (26) and is plotted in Fig. 2. As noted above in Eq. (27), its frequency integral gives the steady-state photonic occupation and it can be shown to diverge at the Dicke transition according to Eq. (13). This result will be explicitly derived in Sec. V A using a low-frequency effective description of  $G^K$ . The photon number diverges in steady state despite the fact that the system undergoes photon loss, and no explicit photon pumping occurs within the model. The reason is that the coupling constant  $g$  is an effective parameter, which in any concrete physical realization microscopically involves a coherent

laser drive process compensating for the loss.

The matrix structure of the Keldysh Green's function  $G_{2 \times 2}^K(\omega)$  can be conveniently exploited to compute the quadrature fluctuations for a general phase angle  $\theta$

$$\langle x_\theta(\omega)x_\theta(-\omega) \rangle = \frac{i}{4\omega_0} \begin{pmatrix} e^{-i\theta} & e^{i\theta} \end{pmatrix} G_{2 \times 2}^K(\omega) \begin{pmatrix} e^{i\theta} \\ e^{-i\theta} \end{pmatrix}, \quad (50)$$

where  $x_\theta$  is defined by

$$x_\theta = \frac{1}{\sqrt{2}\omega_0} (e^{i\theta}a + e^{-i\theta}a^\dagger). \quad (51)$$

The corresponding equal time, frequency-integrated fluctuations  $\langle x_\theta(t)x_\theta(t) \rangle = \int_\omega \langle x_\theta(\omega)x_\theta(-\omega) \rangle$  are plotted in Fig. 3 and diverge at the Dicke transition for all angles  $\theta$  except for  $\theta^*$  defined by  $\theta^* = \pi - \tan^{-1}(\omega_0/\kappa)$ . The angle  $\theta^*$  can also be obtained as the phase angle for the non-diverging eigenmode of the zero-frequency limit of the Keldysh correlation function, thereby naturally yielding the attenuated and amplified quadratures alluded to in Sec. II.

Here in the case of the driven Dicke model, the *equal-time photon fluctuations* of  $x_{\theta^*}$  inside the cavity are independent of the atom-photon coupling and in particular are not attenuated below the vacuum noise level (the  $g = 0$  limit, without atoms in the cavity). This is different from the case of the optical parametric oscillator [43] where, at threshold, the equal-time fluctuations of the non-diverging intra-cavity quadrature are reduced to 50% of the vacuum level. This difference can be traced back the different frequency dependencies of the effective driving term in both situations. In the parametric oscillator, the driving of the cavity occurs via a classically-treated photon pump laser, and the driving amplitude is typically set to a constant coherent field amplitude. In the present case of the driven Dicke model, the effective driving of the cavity is

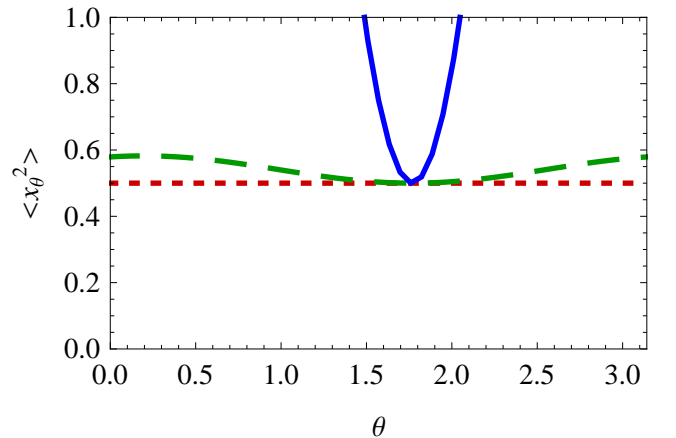


FIG. 3: Frequency integrated equal-time correlations  $\langle (x_\theta(t))^2 \rangle$  as function of the angle  $\theta$ , for different values of the coupling strength approaching the transition at  $g_c \approx 0.51$ :  $g=0$  (dotted), 0.25 (dashed), 0.5 (solid). The fluctuations in the quadrature  $\theta^* = \pi - \tan^{-1}(\omega_0/\kappa) \approx 1.768$  are independent of the coupling strength. Numerical parameters:  $\omega_0 = \omega_z = 1.0$ ,  $\kappa = 0.2$



mediated by the atoms via virtual absorption and emission of photons from the pump laser into the cavity. The corresponding driving term  $\sim g^2 \omega_z / (\omega_z^2 - \omega^2)$  is maximal for frequencies  $\omega$  of the order of the atomic detuning  $\omega_z$  and vanishes for large frequencies.

Nevertheless, also in the driven Dicke model, the experimentally relevant homodyne spectrum  $G_{2 \times 2}^{\text{out}}(\omega)$  of the cavity output field shows noise reduction below the vacuum level in the  $\theta^*$  quadrature [4]. Following the standard “input-output theory” [43, 44], it is possible to show that the homodyne spectrum can be linked to the Keldysh response function via

$$i G_{2 \times 2}^{\text{out}}(\omega) = \left\| \begin{pmatrix} 2i\kappa & 0 \\ 0 & -2i\kappa \end{pmatrix} G_{2 \times 2}^R(\omega) - \mathbf{1}_{2 \times 2} \right\|^2 - \mathbf{1}_{2 \times 2}, \quad (52)$$

where  $\|M\|^2 \equiv M^\dagger M$ . By applying the transformation (50) to  $G_{2 \times 2}^{\text{out}}$ , it is then possible to compute the fluctuations of the output quadrature  $\langle x_{\text{out},\theta}(\omega) x_{\text{out},\theta}(-\omega) \rangle$ . This quantity has been studied in detail in Ref. [4]: at the Dicke transition, the zero-frequency component of  $\theta = \theta^*$  tends to the maximally attenuated value of  $-1/(4\omega_0)$ . It should be noted, however, that the equal-time, frequency integrated fluctuations  $\langle x_{\text{out},\theta^*}(t) x_{\text{out},\theta^*}(t) \rangle$  are zero and therefore not attenuated below vacuum level.

### C. Comparison with a closed system at equilibrium

In the case of a closed system at equilibrium, the retarded Green’s function is the same as Eq. (45), up to the replacement  $\kappa \rightarrow \epsilon$ , and letting  $\epsilon \rightarrow 0$ . The corresponding spectral response function  $\mathcal{A}_{aa^\dagger}(\omega)$  is similar to the one shown for the Markovian case, with the narrow peaks substituted by delta-functions at the resonant frequencies. The Keldysh component of the inverse Green’s function reads, at zero temperature:

$$D_{2 \times 2, \text{EQ}}^K = \begin{pmatrix} 2i\epsilon \text{sign}(\omega) & 0 \\ 0 & -2i\epsilon \text{sign}(\omega) \end{pmatrix}. \quad (53)$$

Again note the different symmetry under frequency reflection with respect to (49).

### D. Photon distribution function and low-frequency effective temperature

We now show that the above mentioned difference between the Markovian case and the equilibrium case leads to the generation of a “low frequency effective temperature” (LET) for the former. For this purpose, we calculate the distribution matrix  $F$ , defined in Eq. (30). To this end, recall that at thermal equilibrium,  $F = \coth(\omega/2T)\mathbf{1}$ : it exponentially approaches unity at high frequencies ( $|\omega| \gg T$ ), and diverges as  $2T/\omega$  at low frequencies ( $|\omega| \ll T$ ). For our Markovian problem, using

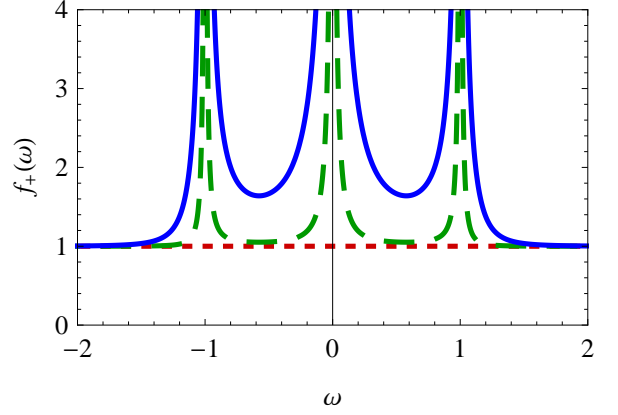


FIG. 4: (Color online) Positive eigenvalue of the distribution function  $F(\omega)$  for the same parameters as in Fig. 2. At low frequencies the distribution diverges as  $2T_{\text{eff}}/\omega$ , where the finite effective temperature is proportional to the photon-atom interaction,  $T_{\text{eff}} \sim g^2$ .

Eq. (43) we find:

$$F = \sigma_z + \frac{2}{\omega} \frac{g^2 \omega_z}{\omega^2 - \omega_z^2} \sigma_x, \quad (54)$$

where  $\sigma_z$  and  $\sigma_x$  are Pauli matrices. The  $F$  matrix is hermitian and traceless, so its two eigenvalues are real and opposite:

$$f_{\pm}(\omega) = \pm \sqrt{1 + \left( \frac{2g^2/\omega_z}{\omega} \frac{1}{1 - (\omega/\omega_z)^2} \right)^2}. \quad (55)$$

Fig. 4 shows the behavior of the positive eigenvalue in two points of the phase diagram, nearby and far away from the transition. In both cases, at low frequencies the eigenvalue diverges as  $1/\omega$ . Exploiting the analogy to the equilibrium case, we obtain the LET

$$T_{\text{eff}} = \frac{g^2}{\omega_z}. \quad (56)$$

We find that  $T_{\text{eff}}$  is not proportional to the decay rate  $\kappa$ . The temperature is rather proportional to the effective interaction between the spin and the photon.  $g$  is the scale that leads to a competition of unitary and dissipative dynamics, and the LET is a measure of this: For  $g = 0$ , the steady state of the dissipative part of the dynamics (empty cavity) is an eigenstate of the Hamiltonian, while this is no longer the case for any finite  $g$ . In the limiting case of  $g \rightarrow 0$ , the LET goes to zero.

A closer inspection of (55) reveals an important difference between the Markovian bath and thermal equilibrium, related to the presence of a second energy scale,  $\omega_z$ . If  $\omega_z \gg T_{\text{eff}}$ , this energy scale does not affect the crossover between the quantum and classical regimes, which then proceeds monotonously similar to an equilibrium problem. If, on the other hand,  $\omega_z \ll T_{\text{eff}}$ , the quantum-classical crossover of the Markovian bath occurs in an unusual way, highlighting the non-equilibrium nature of the problem. Starting from a divergence at zero frequency, the distribution function (55) decreases as

$T_{\text{eff}}/\omega$ , in analogy to an equilibrium system at finite temperature. Then, instead of monotonously decreasing towards the quantum regime where  $f \approx 1$ , it exhibits a second divergence at  $\omega = \omega_z$ . Since the spectral weight vanishes sufficiently fast in this regime, the correlation functions still remain finite; the pole in  $F$  accounts for a different scaling of correlations and spectral properties in this regime (cf. Fig. 2). At higher frequencies, it finally tends to one, following the non-equilibrium curve  $f \approx T_{\text{eff}}/(\omega - \omega_z)$ . The approach to the quantum regime  $f \approx 1$  is polynomial, unlike the exponential approach in the equilibrium case. In App. D, we show that the thermal  $1/\omega$  divergence, leading to a finite LET, is generic for Markovian systems.

We note that our definition of an effective temperature is not the one commonly used in the context of laser theory [24]. In this context, the effective temperature is used only to describe the fluctuations of the photonic field, as compared to the equilibrium fluctuations in the lab frame. As a consequence, the divergence of the photon number at the phase transition is always associated to a diverging effective temperature. In contrast, our low-frequency effective temperature (LET) describes the ratio between the fluctuations and the response of the system in the rotating frame. It is finite at the transition and, as we will see, allows us to map the Dicke transition to an existing dynamical universality class of equilibrium systems.

## V. PHOTON LANGEVIN EQUATION FROM KELDYSH ACTION AND FINITE- $N$ CORRECTIONS

We now rederive the simple scalar Langevin equation (7) as the stochastic equation of motion of the Keldysh action. We also determine how the photon correlations scale with system size  $1/N$  when approaching the Dicke transition.

### A. Low-frequency effective Langevin equation for $N \rightarrow \infty$

In this subsection we derive a simple description of the photon-only action (43), focusing on the  $x = (a + a^*)/\sqrt{2\omega_0}$  quadrature, by mapping its Keldysh action to a linear stochastic equation. The resulting Langevin equation coincides with the equation of motion of a classical particle in an harmonic confinement, coupled to an equilibrium bath at finite temperature consistent with the Heisenberg-Langevin analysis of Sec. II A.

To obtain the Keldysh action of  $x$  we replace  $a = \sqrt{\omega_0/2}(x + ip)$  and  $a^* = \sqrt{\omega_0/2}(x - ip)$  inside Eq. (43) and integrate out the  $p$ -component. This procedure gives:

$$S_{xx} = \int_{\omega} (x_{cl}(-\omega) x_q(-\omega)) \begin{pmatrix} 0 & [G_{xx}^A(\omega)]^{-1} \\ [G_{xx}^R(\omega)]^{-1} & D_{xx}^K(\omega) \end{pmatrix} \begin{pmatrix} x_{cl}(\omega) \\ x_q(\omega) \end{pmatrix}, \quad (57)$$

where

$$D_{xx}^K(\omega) = 2i\kappa \frac{\omega_0^2 + \kappa^2 + \omega^2}{\omega_0}, \quad [G_{xx}^R(\omega)]^{-1} = \frac{1}{\det[G_{2 \times 2}^R(\omega)]}, \quad (58)$$

and the determinant of  $G_{2 \times 2}^R(\omega)$  is given by Eq. (48).

As well known, any quadratic Keldysh action is equivalent to a linear Langevin equation. Starting from a generic quadratic action (57) one introduces a Hubbard-Stratonovich “noise” field  $f(\omega)$  to obtain:

$$S_{xf} = \int_{\omega} 2 \{ [G_{xx}^R(\omega)]^{-1} x_{cl}(\omega) - f(\omega) \} x_q(-\omega) - \frac{f(-\omega)f(\omega)}{D_{xx}^K(\omega)}. \quad (59)$$

(This action is equivalent to Eq. (57), as can be explicitly shown by performing the Gaussian integral over  $f(\omega)$ , and using  $G^R(\omega) = G^A(-\omega)$  – see Ref. 21, 22 for more details.) Inside the Keldysh partition function  $Z^K = \int D\{x_q; x_{cl}; f\} e^{iS_{xf}}$ , the integration over  $x_q$  then takes the form of a delta function with argument

$$[G_{xx}^R(\omega)]^{-1} x_{cl}(\omega) = f(\omega). \quad (60)$$

The remaining last part of the action Eq. (59) involves only  $f(\omega)$  and can be thought of as the statistical weight of a Gaussian random variable with correlations

$$\langle f(\omega)f(\omega') \rangle = -\frac{i}{2} D_{xx}^K(\omega) \delta(\omega + \omega'). \quad (61)$$

Eq. (60) then becomes a stochastic equation of motion for the bosonic field  $x(\omega) = x_{cl}(\omega)$ , identical to the Langevin equation (7) obtained in Sec. II A.

### B. Dynamic critical exponent

In addition to the photon flux exponent, we identify a second indicator of criticality, the dynamical exponent. This exponent governs the decay of the two-time correlations close to criticality. Going back to the Langevin equation (60) and keeping the  $\omega^2$  terms we obtain:

$$(V\omega^2 + 2i\kappa\omega + \alpha^2)x(\omega) = f(\omega), \quad (62)$$

$$\langle f(\omega)f(\omega') \rangle = \kappa \frac{\omega_0^2 + \kappa^2 + \omega^2}{\omega_0} \delta(\omega + \omega'), \quad (63)$$

where we defined a dimensionless parameter  $V = 1 + 4\omega_0 g^2/\omega_z^3$ . For simplicity we further approximate  $V \approx 1$  and obtain the correlation function [55]

$$\begin{aligned} \langle \{\hat{x}(t), \hat{x}(0)\} \rangle &= iG_{xx}^K(t) \\ &= i \int \frac{d\omega}{2\pi} e^{i\omega t} \frac{\kappa(\omega_0^2 + \kappa^2 + \omega^2)}{2\omega_0[(2\kappa\omega)^2 + (\alpha^2 - \omega^2)^2]} \\ &= \frac{e^{-i\kappa t}}{8\omega_0 m^2 \alpha^2} \left[ m^2 (\omega_0^2 + \kappa^2 + \alpha^2) \cos(\sqrt{m^2} t) \right. \\ &\quad \left. + \kappa \sqrt{m^2} (\omega_0^2 + \kappa^2 - \alpha^2) \sin(\sqrt{m^2} t) \right], \end{aligned} \quad (64)$$

where we defined  $m^2 = \alpha^2 - \kappa^2$ . In the vicinity of the transition, for  $0 < \alpha < \kappa$ , the frequency  $m$  becomes purely imaginary and the oscillatory behavior in the above expression disappears.

As already mentioned, this is a generic feature of dissipative phase transitions (see App. C). For sufficiently large times and approaching the transition  $\alpha \rightarrow 0$  (where only the closest pole to zero contributes), we then obtain

$$\langle \{\hat{x}(t), \hat{x}(0)\} \rangle = \frac{\omega_0^2 + \kappa^2}{8\omega_0\alpha^2} e^{-t/\xi_t}. \quad (65)$$

The correlation time:

$$\xi_t = \frac{2\kappa}{\alpha^2} \sim \frac{1}{|g_c - g|^{\nu_t}} \quad (66)$$

is governed by a dynamical exponent  $\nu_t = 1$ .

### C. Finite size scaling

We now include non-linearities for the photon dynamics arising from a finite number of atoms  $N$ . The most relevant term in the Langevin equation is a frequency independent cubic term:

$$(2\kappa\partial_t + \alpha^2)x(t) + \beta^3 x^3(t) = f(t). \quad (67)$$

This equation defines the dynamical critical theory of an Ising transition with no conserved quantities, the so-called “Model A”, for  $n = 1$  degrees of freedom and  $d = 0$  dimensions [34]. The frequencies  $\alpha$  and  $\beta$  need to be determined from the microscopic theory. The former has already been derived in Eq. (9). To obtain  $\beta$ , on the other hand, we demand that the effective Langevin equation reproduces the same saddle-point as the original action Eq. (35). Using  $x = (a + a^*)/\sqrt{2\omega_0}$  and Eq. (38) we then obtain:

$$\sqrt{\frac{-\alpha^2}{\beta^3}} = \sqrt{N} \frac{\sqrt{2\omega_0(g^2 - g_c^2)}}{\kappa^2 + \omega_0^2} \Rightarrow \beta^3 = 2 \frac{(\kappa^2 + \omega_0^2)^2}{N\omega_z}. \quad (68)$$

The parameter  $\beta$  describes deviations from the mean-field solution, due to the finiteness of the system and, indeed, it vanishes in the thermodynamic limit  $N \rightarrow \infty$ .

With Eq. (13) and  $\langle x^2 \rangle = \frac{\int dx x^2 e^{-F(x)/T_{\text{eff}}}}{\int dx e^{-F(x)/T_{\text{eff}}}}$  and the effective free energy  $F(x) = \frac{1}{2}\alpha^2 x^2 + \frac{1}{4}\beta^3 x^4$ , we can write the photon occupation at the critical point as:

$$2\langle n \rangle_c + 1 = 2\omega_0 \left( 1 + \frac{\kappa^2}{\omega_0^2} \right) \sqrt{\frac{T_{\text{eff}}}{\beta^3}} \frac{\Gamma(3/4)}{\Gamma(1/4)} \sim \sqrt{N} \quad (69)$$

Here  $\Gamma(3/4)/\Gamma(1/4) \approx 0.338$  is the ratio of two Gamma functions. We find that the photon occupation at the critical point scales like  $N^{1/2}$ . This result indicates a further distinction between the present model and the zero-temperature quantum phase transition [45, 46], for which  $\langle n \rangle \sim N^{1/3}$ .

To substantiate our findings we compare them with the numerical simulation of the Dicke model (1) with cavity loss (3) using the Monte Carlo Wave-Function (MCWF) method [47], as implemented in the open-source C++QED library [48]. The numerical results in Fig. 5 compare well to the effective

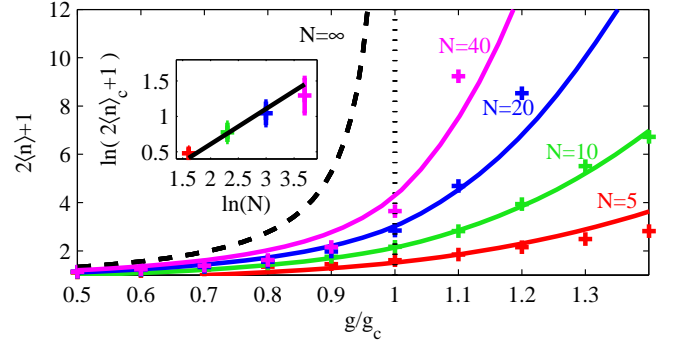


FIG. 5: (Color online) Photon occupation in the vicinity of the Dicke transition for  $N=5$  (red), 10 (green), 20 (blue), 40 (magenta). Crosses: result of the numerical simulation of the full Master equation. Dashed line: mean field result (14), valid in the thermodynamic limit  $N \rightarrow \infty$ . Solid lines: effective Langevin equation, including leading  $1/N$  corrections. The effective model is valid only around the critical point and its validity region narrows with increasing  $N$ . INSET: log-log plot of the photon occupation at the critical point. Crosses: numerical results. Solid line: effective model (69), predicting a  $N^{1/2}$  dependence. Model parameters:  $\omega_z = 2.0$ ,  $\omega_0 = 1.0$ ,  $\kappa = 1.0$ , giving  $g_c = 1.0$ . Other parameters used: number of trajectories  $N_{\text{traj}} = 10$ , time step  $dt = 1$ , total time  $T = 400$ .

Langvin approach (69) and approach the predictions of the mean field approach (14) for  $N \rightarrow \infty$ . We emphasize that no fitting parameters were used when comparing the numerics with the effective model. Independent calculations [49] indicate that the critical exponent of the photon occupation may differ from the predicted rational scaling dimension. A downshift from our analytical result 0.5 to  $0.41 \pm 0.02$  is found. This value is compatible with our present numerical results as well. This discrepancy is probably due to fact that we extract the critical exponent at leading order in the  $1/N$  expansion, which may be a hint at anomalous critical behavior. Such an anomalous dimension can occur at equilibrium as well (both at zero and finite temperature) and does not probe *a priori* the non-equilibrium nature of the transition. However, as the limit  $N \rightarrow \infty$  leads to an approach of a Gaussian fixed point, while any finite value of  $N$  corresponds to a non-critical system, such a result is surprising. A more detailed investigation is left for the future.

We now consider the dynamical properties (i.e. non-equal-time correlations) of the open Dicke model at the critical point and show that they differ from those of an equilibrium system at finite temperature [8–10]. To that end, we study the Langevin equation (67) at the critical point  $\alpha = 0$  under the scaling transformation  $\beta \rightarrow N^{1/3}\beta$  (cf. Eq. (68)). How should we rescale  $t$ ,  $x$ , and  $f$  such that the Langevin equation (67) is still satisfied? First we note that, due to the presence of a finite (effective) temperature  $\langle f(t)f(t') \rangle = T_{\text{eff}}\delta(t - t')$ ,  $f$  scales as  $[t]^{-1/2}$ . The scaling dimensions of  $x$  and  $t$  can be determined by requiring the scale invariance of  $[\partial_t x/f] = [x]/[t]^{1/2}$  and  $[\beta x^3/f] = [x]^3/(N[t]^{-1/2})$ . The resulting scaling of  $[x] = N^{1/4}$  confirms our predictions based on the thermodynamic approach (69). The scaling dimension of the time coordinate  $[t] = N^{-1/2}$  tells us that the spectrum is a scaling function of

$N^{1/2}\omega$ . The present scaling analysis may also explain the recent numerical results of Ref. [49], where the eigen-frequency of the system at the critical point was found to scale as  $1/N^\epsilon$  with  $\epsilon = 0.44 \pm 0.02$ . It should be stressed that, remarkably, our scaling result differs from the one predicted for the finite temperature transition of a closed system [56], for which  $[t] = N^{-1/4}$  and the spectrum is a scaling function of  $N^{1/4}\omega$ .

## VI. KELDYSH APPROACH FOR ATOM OBSERVABLES

We now analyze the single-atom observables of the open Dicke model. To this end, we represent each atom by a real field variable  $\phi_\ell$ , with  $\ell = 1, \dots, N$ . In order for  $\phi_\ell$  to be a faithful representation of an Ising spin, we constrain its fluctuations to lay on a unit Bloch sphere. This single-atom representation of the Ising spins allows to treat the qualitative effects of atom decay within a simplified “friction model” for  $\phi_\ell$  which we explain below. This process couples directly to the local Ising degrees of freedom of the single atoms. The same is true for disorder due to spatial variations of the qubit-photon couplings [12]. In such cases, one cannot employ the single large- $N$  spin representation of the Dicke model.

We proceed by introducing into the path integral  $N$  Lagrange multipliers  $\lambda_\ell$ , corresponding to a suitable Fourier representation of the delta function [12, 50],  $\delta(\phi_\ell^2 - 1) = \int d\lambda \exp^{i\lambda(\phi_\ell^2 - 1)}$ . On the closed time contour, this amounts to adding the following expression to the action:

$$S_{\lambda,\pm} = \frac{-1}{2\omega_z} \int_t \sum_{\ell=1}^N \left[ \lambda_{\ell,+}(t) (\phi_{\ell,+}^2(t) - 1) - \lambda_{\ell,-}(t) (\phi_{\ell,-}^2(t) - 1) \right]. \quad (70)$$

Moving to the “classical/quantum” notation and adding the bare action of the atoms we obtain:

$$S_{\phi,\lambda} = \frac{1}{\omega_z} \int_\omega \sum_{\ell=1}^N (\phi_{\text{cl},\ell}(-\omega) \phi_{\text{q},\ell}(-\omega)) G_{\phi,\lambda}^{-1} \begin{pmatrix} \phi_{\text{cl},\ell}(\omega) \\ \phi_{\text{q},\ell}(\omega) \end{pmatrix} + \frac{1}{\omega_z} \int_t \sum_{\ell=1}^N \lambda_{\text{q},\ell}(t), \quad (71)$$

$$G_{\phi,\lambda}^{-1} = \begin{pmatrix} -\lambda_{\text{q},\ell} & \omega^2 - \lambda_{\text{cl},\ell} + \Sigma_{\phi,\ell}^{\text{A}}(\omega) \\ \omega^2 - \lambda_{\text{cl},\ell} + \Sigma_{\phi,\ell}^{\text{R}}(\omega) & -\lambda_{\text{q},\ell} + \Sigma_{\phi,\ell}^{\text{K}}(\omega) \end{pmatrix}.$$

Note that  $\sqrt{\lambda_{\text{cl},\ell}}$  can be associated with the excitation energy of the Ising spins and is to be determined self-consistently. The atom self-energies  $\Sigma_{\phi,\ell}^{\text{R/A/K}}(\omega)$  will be explained below.

Finally, we have for the atom-cavity interaction:

$$S_{\phi a} = \int_t \sum_{\ell=1}^N \frac{g}{2} [\phi_{+,\ell}(t) (a_+(t) + a_+^*(t)) - \phi_{-,\ell}(t) (a_-(t) + a_-^*(t))]. \quad (72)$$

The Keldysh action for the full Dicke model (1) then becomes

$$S[a, \phi, \lambda] = S_a + S_{\phi,\lambda} + S_{\phi a}, \quad (73)$$

with the various terms given by Eqs. (17, 71, 72).

We model local, single-atom damping in a simple effective way which is consistent with symmetry properties of our real-valued Ising oscillators  $\phi$ . The atoms are subject to decay into photon modes outside the cavity and possible other damping mechanisms like s-wave scattering with other momentum-modes, trap loss or finite-size dephasing [51]. As a result some fraction of the atoms leave the two-density mode Hilbert space which maps to the Dicke model; others may be spontaneously scattered back in. Representing the atoms by a complex field  $\Phi^*, \Phi$ , we subsume the above processes into Markovian decay of the atoms with the self-energies

$$\Sigma_{\Phi^*}^{\text{A}} = -i\gamma, \quad \Sigma_{\Phi^*}^{\text{R}} = +i\gamma, \quad \Sigma_{\Phi^*}^{\text{K}} = 2i\gamma, \quad (74)$$

with  $\gamma$  an effective single-atom decay rate. Our effective real-valued Ising field in Eq. (71) may be viewed as the real component of the originally complex boson  $\Phi_{\text{q/cl}}(t) = \sqrt{\frac{1}{2}} (\phi_{\text{q/cl}}(t) + i\tilde{\phi}_{\text{q/cl}}(t))$ . Integrating out the  $\tilde{\phi}$ -component,

$$\Sigma_{\phi}^{\text{A}}(\omega) = -i\gamma\omega, \quad \Sigma_{\phi}^{\text{R}}(\omega) = +i\gamma\omega, \quad \Sigma_{\phi}^{\text{K}}(\omega) = i\gamma \frac{\omega^2 + \gamma^2 + \omega_z^2}{2\omega_z}. \quad (75)$$

Note that this simple model for spontaneous emission couples to the  $\sigma_x$  projection of the atomic states and does not specify the states of the  $\sigma_z$  projection of the spins. We emphasize, however, that the form of the dissipative self-energies is dictated by the combination of low-frequency expansion and the real-valued nature of the Ising field  $\phi$ . In particular, a frequency independent term is ruled out for  $\Sigma^{\text{A,R}}$ . The above results for Markovian baths should be compared with the results for atoms in equilibrium,

$$\Sigma_{\phi,\text{EQ}}^{\text{A}}(\omega) = -i\epsilon\omega, \quad \Sigma_{\phi,\text{EQ}}^{\text{R}}(\omega) = +i\epsilon\omega, \quad \Sigma_{\phi,\text{EQ}}^{\text{K}}(\omega) = 2i\epsilon\omega \coth\left[\frac{\omega}{2T}\right], \quad (76)$$

where one also lets  $\epsilon \rightarrow 0$  at the end of the calculation. We here analyze the Dicke model in terms of the atomic degrees of freedom alone. One can exactly integrate out the photons from the action (73). This is conveniently done by going to a coordinate representation of the photons:  $a_{\text{q/cl}}(t) = \sqrt{\frac{\omega_0}{2}} (x_{\text{q/cl}}(t) + ip_{\text{q/cl}}(t))$ , and first performing the integration over  $p_{\text{q/cl}}$  and subsequently over  $x_{\text{q/cl}}$ . We obtain, with  $S_{\phi,\lambda}$  given by Eq. (71), the atom-only action,

$$S[\phi, \lambda] = S_{\phi,\lambda} + S_{\phi,\text{g}^2}, \quad (77)$$

$$S_{\phi,\text{g}^2} = -\frac{1}{2} \int_\omega \sum_{\ell,m=1}^N \frac{g^2}{N} \times \quad (78)$$

$$(\phi_{\text{cl},\ell}(-\omega) \phi_{\text{q},\ell}(-\omega)) \begin{pmatrix} 0 & \sigma^{\text{A}}(\omega) \\ \sigma^{\text{R}}(\omega) & \sigma^{\text{K}}(\omega) \end{pmatrix} \begin{pmatrix} \phi_{\text{cl},m}(\omega) \\ \phi_{\text{q},m}(\omega) \end{pmatrix},$$

where the matrix entries are

$$\begin{aligned}\sigma^R(\omega) &= [\sigma^A(\omega)]^* = \frac{-2\omega_0}{(\omega + i\kappa)^2 - \omega_0^2}, \\ \sigma^K(\omega) &= \frac{2i\kappa(\omega^2 + \kappa^2 + \omega_0^2)}{[(\omega - i\kappa)^2 - \omega_0^2]^2}.\end{aligned}\quad (79)$$

We now decouple Eq. (78) with a Hubbard-Stratonovich field  $\psi_\ell(\omega) \leftrightarrow \phi_\ell(\omega)$  and integrate out the  $\phi$ -field. We assume  $\psi$  to

be time-independent,  $\psi(\omega) \equiv \psi/(2\pi)\delta_{\omega,0}$ , and substitute the  $N$  Lagrange multipliers by a single effective field  $\lambda_{q/cl,\ell} \rightarrow \lambda_{q/cl}$ . The resulting Keldysh partition function

$$Z^K = \int D\psi D\lambda e^{i\frac{N}{2\pi}S[\psi,\lambda]}, \quad (80)$$

obtains a prefactor of the number of atoms  $N$  in the exponent multiplying the action

$$\begin{aligned}S[\psi, \lambda] &= \frac{\lambda_q}{\omega_z} + g^2 \left( \sigma^R(0)\psi_{cl}\psi_q + \frac{1}{2}\sigma^K(0)\psi_q^2 \right) + \frac{i}{2} \int_\omega \left( \ln \left[ \lambda_q \left( \Sigma_\phi^K(\omega) - \lambda_q \right) + \left( \omega^2 - \lambda_{cl} + \Sigma_\phi^A(\omega) \right) \left( \omega^2 - \lambda_{cl} + \Sigma_\phi^R(\omega) \right) \right] \right) \\ &\quad - \frac{g^4}{4} \frac{\omega_z}{\lambda_q \left( \Sigma_\phi^K(0) - \lambda_q \right) + \lambda_{cl}^2} \left[ \left( \Sigma_\phi^K(0) - \lambda_q \right) \left( \sigma^R(0) \right)^2 \psi_q^2 + 2\lambda_{cl}\sigma^R(0)\psi_q \left( \sigma^R(0)\psi_{cl} + \sigma^K(0)\psi_q \right) - \lambda_q \left( \sigma^R(0)\psi_{cl} + \sigma^K(0)\psi_q \right)^2 \right].\end{aligned}\quad (81)$$

Taking  $N \rightarrow \infty$ , we now extract the phase diagram, response and correlation functions, and the value of the order parameter using a saddle-point approximation. This can be obtained by requiring the derivatives with respect to  $\lambda_q$  and  $\psi_q$  to be zero, and then substituting  $\lambda_q = 0$ ,  $\lambda_{cl} = \lambda$ ,  $\psi_q = 0$ ,  $\psi_{cl} = \psi$ . The derivative with respect to  $\lambda_q$  constrains –by construction Eq. (70)– the frequency integral of the Keldysh Green's function to be equal to unity:

$$\frac{\partial S}{\partial \lambda_q} = 0 \quad \Rightarrow \quad \langle \phi^2 \rangle = \int \frac{d\omega}{2\pi} iG_{\phi\phi}^K(\omega) = 1 \quad (82)$$

with

$$\begin{aligned}G_{\phi\phi}^K(\omega) &= \frac{-\omega_z \Sigma_\phi^K(\omega)}{2 \left( \omega^2 - \lambda + \Sigma_\phi^A(\omega) \right) \left( \omega^2 - \lambda + \Sigma_\phi^R(\omega) \right)} \\ &\quad - 2\pi i \delta(\omega) \frac{g^2 \omega_z^2}{4\lambda^2} \left( \sigma^R(0) \right)^2 \psi^2.\end{aligned}\quad (83)$$

The saddle-point condition for the order parameter yields

$$\frac{\partial S}{\partial \psi_q} = 0 \quad \Rightarrow \quad \psi \left[ 1 - \frac{g^2 \omega_z}{2} \frac{\sigma^R(0)}{\lambda} \right] = 0. \quad (84)$$

To determine the position of the Dicke transition, we need to compute the saddle point value of  $\lambda_{cl} = \lambda$  in the normal and in the superradiant phases and equate the two values. In the normal (N) phase  $\lambda$  is determined by (82) and (83) with  $\psi = 0$ ,

$$\lambda_N = \frac{\gamma^2 + \omega_z^2}{3}. \quad (85)$$

Note that naively taking  $\gamma \rightarrow 0$  does not reproduce the equilibrium value for  $\lambda_N$ , cf. also Subsec. VI C. In the ferromagnetic phase (FM) the order parameter acquires a finite expectation value  $\psi \neq 0$  and, to fulfill Eq. (84), we need to require the

argument of the square bracket to be zero:

$$\lambda_{FM} = \frac{g^2 \omega_0 \omega_z}{\omega_0^2 + \kappa^2}, \quad (86)$$

where we have used Eq. (79) for  $\sigma^R(0)$ . At the phase boundaries both (85) and (86) must hold, leading to:

$$g_c = \sqrt{\frac{(\gamma^2 + \omega_z^2)(\kappa^2 + \omega_0^2)}{3\omega_0 \omega_z}}. \quad (87)$$

Eqs. (82) and (84) determine the value of ferromagnetic order parameter  $\psi$  as well. When approaching the phase transition from above ( $g \geq g_c$ ),  $\psi$  vanishes as

$$\psi = \sqrt{\frac{3(g + g_c)}{4g^2}} \sqrt{g - g_c}. \quad (88)$$

Compared with a closed system at equilibrium, Eq. (92), the order parameter for this open Markovian system vanishes with enhanced amplitude but with the same mean-field like square-root exponent.

#### A. Atom spectral response function

The single-atom retarded and advanced Green's function are determined by the derivative of (81) with respect to  $\lambda_{cl}$  and reads

$$G_{\phi\phi}^R(\omega) = [G_{\phi\phi}^A(\omega)]^* = \frac{\omega_z}{2 \left( \omega^2 - \lambda + \Sigma_\phi^R(\omega) \right)}, \quad (89)$$

from which follows the spectral response function

$$\mathcal{A}_{\phi\phi}(\omega) = -2\text{Im}G_{\phi\phi}^R(\omega) \quad (90)$$

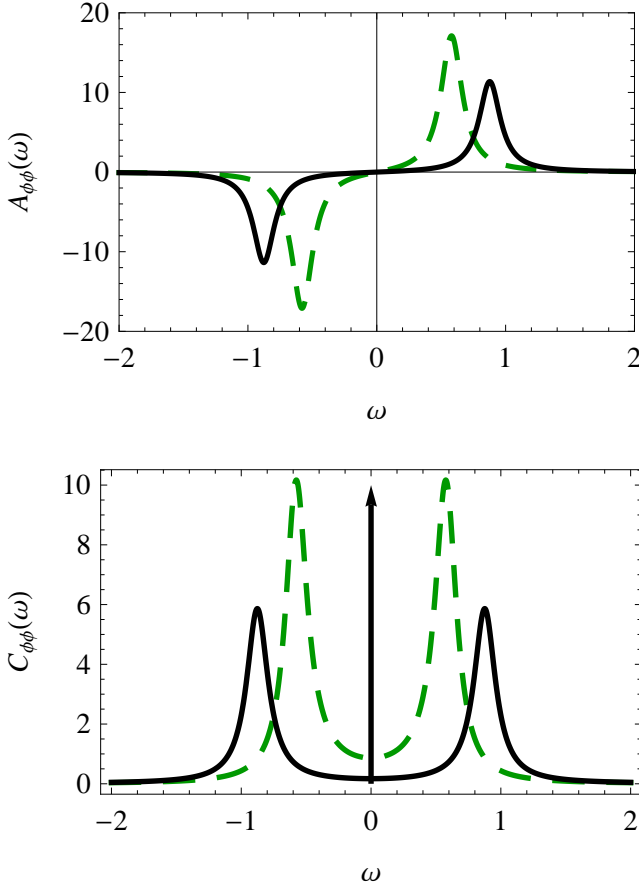


FIG. 6: (Color online) Atomic spectral response function  $\mathcal{A}_{\phi\phi}(\omega)$  and atomic correlation function  $C_{\phi\phi}(\omega)$  in the normal phase ( $g < g_c$ , green, dashed) and in the superradiant phase ( $g = 1.2g_c$ , black, solid). Other parameters used:  $\omega_0 = \omega_z = 1$ ,  $\gamma = \kappa = 0.2$  (leading to  $g_c \approx 0.6$ ). The arrow illustrates the delta-function contributions from Eq. (83) in the superradiant phase.

as the expected frequency-resolved signal from the atoms after local, time-modulated density perturbations. Fig. 6 displays the characteristic Lorentzian shape of the spectral response function peaked at  $\sqrt{\lambda}$ , broadened by single atom decay  $\sim \gamma$ . The single atom response is smooth across the transition. This is not to be confused with the roton-type mode softening observed by Mottl *et al.* [51] that pertains to the *collective* atomic density excitations for a finite number of atoms.

### B. Atom correlation function

The atom correlation function is given by the Keldysh Greens function Eq. (83)

$$C_{\phi\phi}(\omega) = iG_{\phi\phi}^K(\omega), \quad (91)$$

$G_{\phi\phi}^K$  is defined in (83), and is exhibited in Fig. 6. Even before the onset of the superradiance peaks (black and blue-dashed arrows) for  $g \geq g_c$ , the correlation function has finite weight

at  $\omega = 0$ . This is the non-equilibrium signature from spontaneous emission of the pumped atoms into the vacuum outside the cavity (a continuum of modes with characteristic frequencies orders of magnitudes lower than the optical photons the pumped atoms emit when they spontaneously decay).

### C. Comparison with a closed system at equilibrium

Note that taking  $\kappa \rightarrow 0$  and  $\gamma \rightarrow 0$  in Eq. (87) does not reproduce the equilibrium value for a closed system. This is due to non-commuting limits of making the Markov approximation and performing the integral to fulfill the sum rule Eq. (82). To obtain  $g_c^{\text{EQ}}$  one needs to use the equilibrium bath self-energies Eqs. (28,76) from the start of the calculation and obtains the equilibrium analogs of Eqs. (87,88):

$$\begin{aligned} g_c^{\text{EQ}} &= \frac{1}{2} \sqrt{\omega_z \omega_0}, \\ \psi^{\text{EQ}} &= \frac{1}{\sqrt{g}} \sqrt{g - g_c}. \end{aligned} \quad (92)$$

We note that, both dissipative channels, cavity photon loss and atomic spontaneous emission, shift the critical value of the coupling. The amplitude with which the ferromagnetic order parameter vanishes is also different.

Note that in a model for “one-way” spontaneous emission coupling to  $\sigma^+$  and  $\sigma^-$  starting from a fully polarized atomic state as assumed in Sec. II A, one would recover the equilibrium limit for  $\gamma \rightarrow 0$ . As explained above, our “mixing” model for spontaneous emission couples to  $\sigma^x$  and thereby assumes a mixed state of the atoms (similar to a many-body paramagnet).

### D. Atom distribution function and low-frequency effective temperature

We now execute the procedure of subsection III D to calculate the effective temperature of the atoms. With the atom Green’s functions and the simple model for atom decay presented above, one finds

$$F_{\phi\phi}(\omega) = \frac{C_{\phi\phi}(\omega)}{\mathcal{A}_{\phi\phi}(\omega)} = \frac{G_{\phi\phi}^K(\omega)}{G_{\phi\phi}^R(\omega) - G_{\phi\phi}^A(\omega)} = \frac{\omega^2 + \gamma^2 + \omega_z^2}{2\omega_z} \frac{1}{\omega}, \quad (93)$$

leading to the effective temperature

$$T_{\phi}^{\text{eff}} = \frac{\gamma^2 + \omega_z^2}{4\omega_z}, \quad (94)$$

which is independent of the coupling strength to the photons (cf. also Fig. 7). For  $\gamma \ll \omega_z$ , the effective temperature is set by the recoil energy of the atoms  $E_R = \omega_z/2$ . Within our model,  $T_{\phi}^{\text{eff}}$  also does not depend on the cavity loss rate  $\kappa$ , contrary to what obtains for the model of Ref. 13. In our case the reason for this is the careful treatment of the thermodynamic



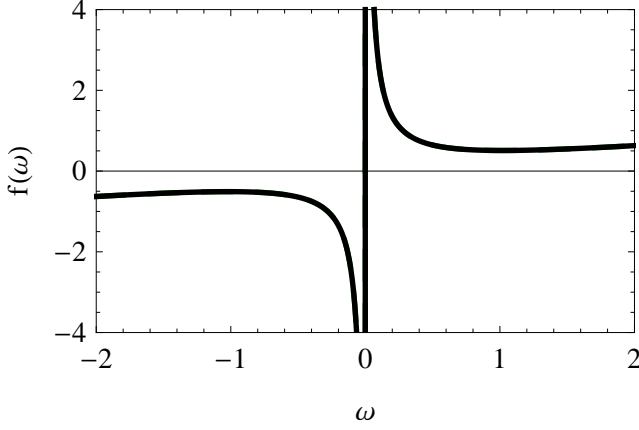


FIG. 7: (Color online) Distribution function  $F_{\phi\phi}(\omega)$  for the same numerical parameters as Fig. 6. The expression for the distribution function Eq. (94) is independent of the atom-photon coupling  $g$ .

limit  $N \rightarrow \infty$  limit leading to Eqs. (77,81). This limit ensures that the only photon-induced self-energies for the atoms occur for zero-frequency quantities (the weight of  $\delta(\omega)$  in Eq. (83)).

## VII. CONCLUSION

In this paper, we presented a path integral approach for the non-equilibrium steady-states of driven quantum systems coupled to Markovian baths, such as ultracold atoms in optical cavities. In the past, these systems have more often been described using a Master equation formalism. We believe that our Keldysh approach allows an easier comparison with other equilibrium and non-equilibrium (classical and quantum) systems. While some of the results presented here are actually new, and not just known results re-phrased in a new approach, the full utility of our approach will become clear then computing thermodynamics and critical properties of large, open systems with spatially fluctuating degrees of freedom such as disorder [12, 14]. In these correlated quantum many-body situations, Master equation approaches are typically limited to relatively small number of atoms and a recipe to compute disorder-averaged quantities does not seem to exist.

We first applied our formalism to the cavity vacuum and subsequently added atomic qubits, interacting with the cavity through a Dicke interaction, and computed the key observables for both, the photons and atoms. In the presence of the atoms, the cavity exhibits a thermodynamic “superradiant” phase transition, whose critical properties we analyzed in detail. They differ from the zero-temperature quantum phase transition of the corresponding closed system due to the presence of an emergent finite “effective temperature” for the relevant low-energy photons, induced by the atom-photon interaction. Due to the non-equilibrium nature of the problem, the effective temperature characterizes only the low-frequency behavior of the system. At higher frequencies, the distribution functions display non-equilibrium and quantum behaviors. The definition of an effective temperature is observable

dependent, leading to different effective temperatures for the atoms versus the photons.

In the future, it will be interesting to apply our approach to dissipative quantum glasses coupled to Markovian (and other) baths such as potentially achievable in multi-mode optical cavities [12–14] or circuit QED [52]. It would also be desirable to obtain a more general classification of conditions under which quantum phase transitions of closed systems are turned into thermal phase transitions by dissipation—and perhaps to find counterexamples by engineered dissipation along the lines of Refs. [38, 39].

## Acknowledgments

We thank M. A. Baranov, F. di Piazza, S. Gopalakrishnan, B. Halperin, E. Kessler, D. Marcos, J. Otterbach, P. Rabl, H. Ritsch, L. Sieberer, S. Sachdev, H. Tureci, and P. Zoller for useful discussions and M. Buchhold, H. Ritsch, D. Marcos and S. Gopalakrishnan for critical comments on the manuscript.

This research was supported by the National Science Foundation (NSF), by the Center for Ultracold Atoms (CUA), by the Multidisciplinary University Research Initiative (MURI), by the Packard Foundation, by the DARPA OLE program, by the DFG under grant Str 1176/1-1, by the Austrian Science Fund (FWF) through SFB FOQUS F4016-N16 and the START grant Y 581-N16, by the European Commission (AQUTE, NAMEQUAM), and by the Institut für Quanteninformation GmbH.

## Appendix A: Self-energy of an open cavity in the rotating frame

In this appendix, we discuss how the coherent drive with a frequency scale  $\omega_p$ , which exceeds all other frequency scales, justifies the form of the Markovian dissipative action Eq. (16), which in particular displays frequency independent terms only which are  $\delta$ -correlated in time and neglect memory effects.

Our starting point is the Hamiltonian of a single boson  $a$ , coupled to a continuum of vacuum fields  $\psi_k$ , via

$$H_0 = \omega_c a^\dagger a + \sum_k \omega_k \psi_k^\dagger \psi_k + g_k (a^\dagger \psi_k + a \psi_k^\dagger). \quad (\text{A1})$$

Here  $g_k$  is the coupling constant between the cavity boson and the external vacuum, and we neglected counter-propagating terms of the form  $a^\dagger \psi_k^\dagger$ . Eq. (A1) is quadratic in  $\psi_k$ , allowing us to analytically integrate-out the vacuum fields and obtain a cavity-only action of the form (17). If we assume that the vacuum fields are kept at an equilibrium temperature  $T_{\text{ext}} = 300\text{K}$ , the corresponding entries are

$$\begin{aligned} [G^R(\omega)]^{-1} &= \omega - \omega_c - \delta\omega + iK(\omega), \\ D^K &= 2iK(\omega) \coth\left(\frac{\omega}{2T_{\text{ext}}}\right). \end{aligned} \quad (\text{A2})$$

Here  $\delta\omega$  corresponds to the Lamb shift and can be absorbed

in a finite renormalization of  $\omega_c$ . The function  $K(\omega) = \sum_k |g_k|^2 \delta(\omega - \omega_k)$  is the spectral density of the vacuum.

The inverse Green's functions (A2) describe the cavity mode in the lab frame. In practice, it is often more convenient to move to a frame rotating with a constant frequency, in our case corresponding to the pump frequency  $\omega_p$ . In this frame, the photons are described by Eq. (A2) with  $\omega \rightarrow \omega_p + \omega$ . Next, we apply the equivalent of the Wigner-Weisskopf approximation, by Taylor expanding the inverse Green's function in small  $\omega$  to zero order. This approximation is justified by the energy scale separation discussed in the text. We obtain

$$[G^R(\omega)]^{-1} = \omega - \omega_c - \delta\omega - \omega_p + iK(\omega_p),$$

$$D^K = 2iK(\omega_p) \coth\left(\frac{\omega_p}{2T_{\text{ext}}}\right). \quad (\text{A3})$$

The factor  $\coth\left(\frac{\omega_p}{2T_{\text{ext}}}\right)$  plays the role of the  $2n + 1$  factor appearing in the finite temperature extension of the master equation (2). To be precise, the two expressions coincide only for  $\omega_c = \omega_p$ . For most experiments the pump frequency is anyway much higher than the external temperature and we can approximate  $\coth\left(\frac{\omega_p}{2T_{\text{ext}}}\right) = 1$ . Under this approximation, Eq. (A3) becomes equivalent to Eqs. (20-21) with

$$\omega_0 = \omega_c + \delta\omega - \omega_p, \quad \kappa = K(\omega_p). \quad (\text{A4})$$

### Appendix B: Photon-only action in the superradiant phase

In the superradiant (SR) phase, the field  $a$  ( $b$ ) is given by the sum of a time-independent component  $a_0$  ( $b_0$ ) and a fluctuating term. The action governing the fluctuating terms can be obtained from Eq. (35) by substituting  $a \rightarrow a_0 + \delta a$  ( $b \rightarrow b_0 + \delta b$ ) (we choose  $a_0, b_0$  real without loss of generality). At  $N \rightarrow \infty$ , we end up with the quadratic action

$$S_{\text{SR}} = \frac{1}{2} \int_{\omega} \delta V_8^\dagger(\omega) \begin{pmatrix} 0 & [G_{4 \times 4}^A]^{-1}(\omega) \\ [G_{4 \times 4}^R]^{-1}(\omega) & D_{4 \times 4}^K \end{pmatrix} \delta V_8(\omega) \quad (\text{B1})$$

with the Green's functions  $[G_{4 \times 4}^R]^{-1}(\omega) = ([G_{4 \times 4}^A]^{-1}(\omega))^\dagger =$

$$\begin{pmatrix} \omega - \omega_0 + i\kappa & 0 & -\bar{g} & -\bar{g} \\ 0 & -\omega - \omega_0 - i\kappa & -\bar{g} & -\bar{g} \\ -\bar{g} & -\bar{g} & \omega - \omega_z - 2\delta\omega_z & -\delta\omega_z \\ -\bar{g} & -\bar{g} & -\delta\omega_z & -\omega - \omega_z - 2\delta\omega_z \end{pmatrix},$$

$$D^K = 2i \text{diag}(\kappa, \kappa, 0, 0). \quad (\text{B2})$$

Here the 8-vector  $\delta V_8(\omega)$  is defined in the analogous way to  $V_8(\omega)$  of Eq. (40) and

$$\bar{g} = g - \frac{3g}{N} b_0^2 \approx \frac{3g^2 + g^2}{4g}, \quad (\text{B3})$$

$$\delta\bar{\omega}_z = \frac{4g}{N} b_0(a_0 + a_0^*) \approx -4\omega_z \frac{g^2 - g_c^2}{g^2}. \quad (\text{B4})$$

We note that the principal change to the spectral response and correlation function in the superradiant phase

$$\mathcal{A}_{aa^\dagger}(\omega) = \mathcal{A}_{\delta a, \delta a^\dagger}(\omega), \quad (\text{B5})$$

$$C_{aa^\dagger}(\omega) = C_{\delta a, \delta a^\dagger}(\omega) + a_0^2 \delta_{\omega, 0} \quad (\text{B6})$$

is the  $\delta$ -function peak at  $\omega = 0$  in the correlation function due to coherent photons ("photon condensate").

### Appendix C: Damped dynamics near the phase transition

We argue based on a systematic low-frequency expansion of the inverse retarded Green's function that the overdamped dynamics observed in the vicinity of the phase transition is generic for systems where a phase transition is driven by a competition within the Hamiltonian sector, while dissipative dynamics acts as a "spectator". To see this, we (i) write the most general form of the inverse retarded Green's function

$$[G_{2 \times 2}^R]^{-1}(\omega) = \begin{pmatrix} p(\omega) & o(\omega) \\ o^*(-\omega) & p^*(-\omega) \end{pmatrix}, \quad (\text{C1})$$

and (ii) use that the phase transition is governed by low frequency behavior and an expansion in powers of the frequency is appropriate,

$$p(\omega) = -\nu + z\omega, \quad o(\omega) = -\mu + y\omega \quad (\text{C2})$$

with complex coefficients and a low frequency spectrum

$$\omega_{\pm} = \frac{i\text{Im}[z^* \nu - y^* \mu] \pm \sqrt{(|z|^2 - |y|^2)\alpha^2 - (\text{Im}[z^* \nu - y^* \mu])^2}}{(|z|^2 - |y|^2)}. \quad (\text{C3})$$

Without dynamic renormalization effects,  $z = 1$  and all other frequency coefficients are zero, so they will be generically much smaller than one (more precisely,  $|\text{Im}[z]|, |\text{Re}[y]|, |\text{Im}[y]| \ll 1$ ), and in particular  $|z|^2 \gg |y|^2$ . (In the large  $N$  open Dicke model, they are exactly zero.) A mass gap, i.e. the scale that characterizes the action at zero frequency, provides a measure of the distance from the phase transition and reads

$$\alpha^2 \equiv \det G_{2 \times 2}^{R-1}(\omega = 0) = |\nu|^2 - |\mu|^2 \geq 0 \quad (\text{C4})$$

(the last inequality must hold for a stable physical system). Approaching the phase transition, this gap shrinks to zero, such that the frequencies must become purely imaginary as a generic feature of a phase transition in the presence of dis-



sipation. Indeed, in a situation where the phase transition is driven by a competition within the Hamiltonian sector of the problem by a quantity  $g$ , the dominant  $g$  dependence is contained in  $\alpha^2(g)$  (more precisely,  $\text{Re}[\nu]$ ,  $\text{Re}[\mu]$ ), while the dissipative scales ( $\text{Im}[\nu]$ ,  $\text{Im}[\mu]$ ) do not strongly depend on  $g$  and remain essentially at their bare, finite values even at the transition point. (In the open Dicke model, only the real parts are modified, while the imaginary parts exactly remain at their bare values.) In such a situation, as  $\alpha(g \rightarrow g_c) \rightarrow 0$ , we may expand the square root in Eq. (C3) in  $\alpha^2$ , identifying that parameter as the distance from the phase transition.

#### Appendix D: $1/\omega$ divergence in Markov distribution functions

We here show that the  $1/\omega$  pole in the photon distribution function at low frequency, and the associated low-frequency effective temperature (LET), is indeed a generic feature of Markovian non-equilibrium systems.

To this end, we consider the low-frequency regime, where in the spirit of a systematic derivative expansion the inverse retarded Green's function of the photon takes the form

$$G^{R-1} = (\omega + i\nu_2)\sigma_z - H, \quad H = \nu_1 \mathbf{1} + \mu_1 \sigma_x + \mu_2 \sigma_y, \quad (\text{D1})$$

where  $\nu_1, \mu_1(\nu_2, \mu_2)$  denote the real (imaginary) part of  $\nu, \mu$ . We set  $z = 1, y = 0$  here, which in principle contribute at  $\mathcal{O}(\omega)$ , and anticipate that this omission will not alter the qualitative results. The hermitean part  $H$  represents Hamiltonian dynamics, the antihermitean  $\sim i\sigma_z$  decay. In a derivative expansion, the most general form of the Keldysh component is

$$D^K = 2i(\kappa_1 \mathbf{1} + \kappa_2 \sigma_x). \quad (\text{D2})$$

Solving the fluctuation-dissipation relation, we obtain

$$F(\omega) = \frac{\kappa_1}{\nu_2} \sigma_z - \frac{1}{\omega} \left[ \frac{\kappa_1}{\nu_2} (\mu_1 \sigma_x + \mu_2 \sigma_y) + \frac{\kappa_2 \mu_1}{\nu_2} \mathbf{1} + \kappa_2 \sigma_y \right]. \quad (\text{D3})$$

Crucially, this confirms the  $1/\omega$  divergence behavior of the distribution function. Allowing for the most general form of  $G^{R-1}$  of the frequency expansion in terms of a finite imaginary part of  $z$  and finite  $y$  only results in subleading corrections: also in this case,  $\lim_{\omega \rightarrow 0} [\omega \cdot F(\omega)] \rightarrow \text{const.}$  Clearly, adding frequency dependent terms to  $D^K$  only leads to subleading corrections in  $F$ . Therefore, the  $1/\omega$  pole at low frequency, and the associated scale generated in this regime, the LET, is a generic feature of Markovian non-equilibrium systems.

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